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ELECTRON SPIN RESONANCE SPECTRA OF THE NH_2 AND ND_2 FREE RADICALS AT 4.2°K*

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Electron spin resonances of NH_2 and ND_2 stabilized in argon at liquid helium temperature have been observed. The identification of the radicals has been established by a detailed analysis of the hyperfine spectra.

The essential features of the apparatus have been described.[1-3] Preliminary experiments in which the products of electric discharges in ammonia were condensed in an argon matrix gave evidence for trapped NH_2 radicals. In addition to a nine-line spectrum which was tentatively assigned to NH_2, we observed strong atomic hydrogen lines, atomic nitrogen lines, and a broad unresolved resonance near the free electron value. When the products of a discharge in ND_3 were condensed in argon, the presence of a strong center deuterium atom line and the broad unresolved resonance line obscured the ND_2 resonances.

In view of these difficulties, an attempt was made to produce the radicals by photolysis of ammonia in argon at 4.2°K. A gaseous mixture of about 0.5% NH_3 or ND_3 and argon was deposited on the liquid-helium-cooled sapphire rod and subsequently photolyzed through a quartz window with 1849-A resonance radiation from a low-pressure mercury lamp.

Figure 1 shows the spectrum obtained by photolysis of NH_3. The spectrum consists of three groups of triplet lines, consistent with the expected spectrum for the NH_2 radical since $I(N^{14})=1$ and $I(H)=1/2$ resulting in nine resonance lines. Because the widths of the resonance lines

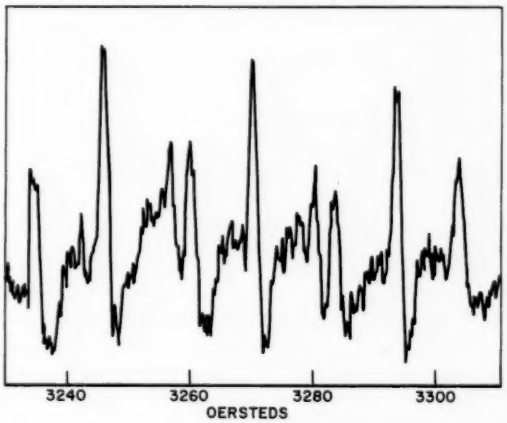

FIG. 1. Electron spin resonance spectrum of NH_2 in an argon matrix at 4.2°K. ν = 9177.82 Mc/sec.

are different, the intensity distributions obtained from peak amplitudes are misleading. In particular, the lines in a triplet group are actually comparable in intensity. It is not evident from the spectrum whether to attribute the splitting into triplet groups to hydrogen and the subsplitting to nitrogen or vice versa. Experiments with deuterated ammonia supplied the answer to this problem and established that the large splitting in NH_2 was due to hydrogen. The spectrum obtained by photolysis of ND_3 in argon is shown in compressed form in Fig. 2. The ND_2 spectrum is expected to consist of 15 lines since $I(N^{14}) = 1$ and $I(D) = 1$. From an analysis of several of the stronger lines of the spectrum, two hyperfine

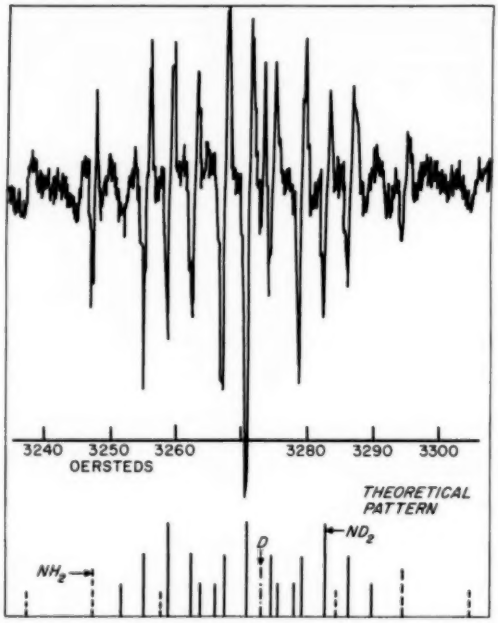

FIG. 2. Electron spin resonance spectrum of ND_2 in an argon matrix at 4.2°K. Also present are spectra of D and NH_2 and weak traces of NHD.

coupling constants were evaluated, as discussed later, and the complete spectrum calculated. The predicted positions of the spectral lines are indicated in the bottom of the figure. Eleven of the lines are clearly recognized in the recording. The others are too close to other lines to be resolved. In addition, one sees the center deuterium atom line, slight traces of NHD, and several lines from NH_2 arising from some remanent NH_3 in the system. The lines of NH_2 and ND_2 were recorded individually on expanded sweeps to determine their field positions with high precision.

If one solves the spin Hamiltonian for the magnetic energy, W, to the first order approximation, one obtains

$$W = M_J g_J \mu_0 H + A M_J \sum_i m_i + B M_J M_I(\mathrm{N}) + \mu_0 H \left[g_I \sum_i m_i + g_I(\mathrm{N}) M_I(\mathrm{N}) \right] \quad (1)$$

where m_i is the nuclear magnetic quantum number of hydrogen ($\pm 1/2$) in the case of NH_2 or of deuterium (1, 0, -1) in the case of ND_2, and the other symbols have their usual significance. The hyperfine coupling constants A and B and the electronic g-factor, g_J, can be calculated from the data. The results are given in Table I.

Table I. Electron spin resonance results for NH_2 and ND_2

Radical	A (Mc/sec)	B (Mc/sec)	g_J
NH_2	67.03(20)	28.90(20)	2.00481 (8)
ND_2	10.27(20)	33.28(20)	2.00466 (8)

The ratio of the hydrogenic coupling constants, $A_H/A_D = 6.526$, is in excellent agreement with the predicted ratio, $g_I(\mathrm{H})/g_I(\mathrm{D}) = 6.514$. The difference between the values for the nitrogen coupling constants is unexpected, indicating that the electronic wave functions for the two radicals are somewhat different. It is clear that the inclusion of higher-order terms in the solution of the spin Hamiltonian would not bring the B values into closer agreement. Apparently, the hyperfine interaction with the nitrogen atom is rather sensitive to some small perturbation in the electronic state. Thus far, we have been unable to account for the discrepancy by considering zero-point vibration and the differences in the rotational states of these molecules.

* This work supported by Bureau of Ordnance, Department of the Navy.

[1] Jen, Foner, Cochran, and Bowers, Phys. Rev. 104, 846 (1956).
[2] Foner, Jen, Cochran, and Bowers, J. Chem. Phys. 28, 351 (1958).
[3] Jen, Foner, Cochran, and Bowers, (to be published).

SPIN EXCHANGE IN SUPERCONDUCTORS

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The only known superconductor among the rare earth elements is lanthanum. The elements following lanthanum in the periodic system are either strongly paramagnetic or ferromagnetic, with magnetic moments which are due to their 4f electrons. In lutetium, 14 electrons have filled this 4f shell entirely and the element does not show pronounced paramagnetism. Lutetium, however, is not superconducting above 1.02°K because its metallic radius has become much smaller and at the same time it is much heavier than lanthanum.

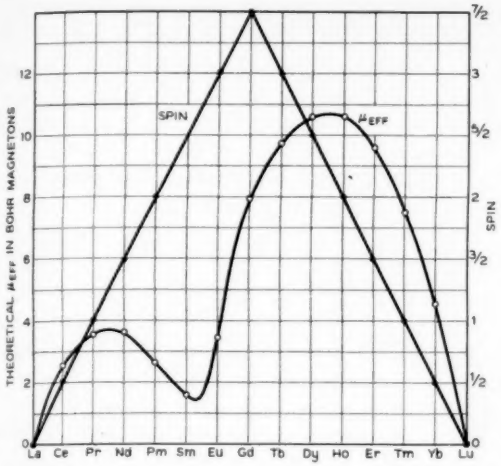

FIG. 1. Effective magnetic moments and spins of the rare earth elements (see reference 2).

The effective magnetic moments of the rare earth elements follow Van Vleck's well-known curve,[1] Fig. 1. These moments, which originate in the low-lying 4*f* shell, are usually assumed to remain undisturbed in almost all chemical compounds which include these elements. It was therefore our hope that by dissolving small amounts of the magnetic rare earth elements in lanthanum, the superconducting transition would be affected by the dipole field from the moment of the rare earth atoms. In Fig. 2 we show the superconducting transitions of lanthanum samples in which 1 at.% of various rare earth elements has been dissolved. It is immediately apparent from these data that the simple assumption that

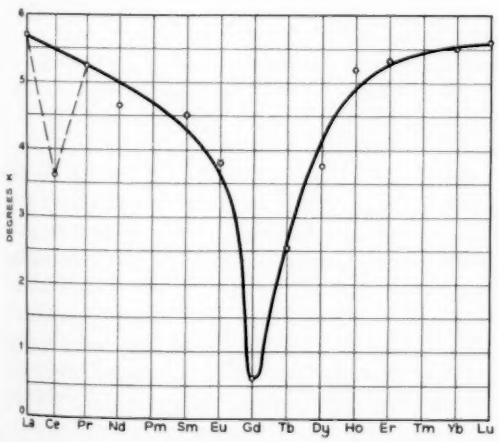

FIG. 2. Superconducting transition temperatures of 1 at.% rare earth solid solutions in lanthanum.

an increase in effective moment should go hand in hand with a decrease of the superconducting transition temperature does not work at all. Instead, the depression of the superconductivity seems to be correlated only with the spin of the solute atoms. (The anomalous behavior of cerium is probably caused by the rather easy shift of part of the 4*f* electron into the 5*d* band; this occurs in the pure metal at low temperature or under pressure.)

The change in the superconducting transition temperature of lanthanum caused by varying the dissolved amounts of gadolinium was investigated in more detail and the results are shown in Fig. 3. The superconducting transition tempera-

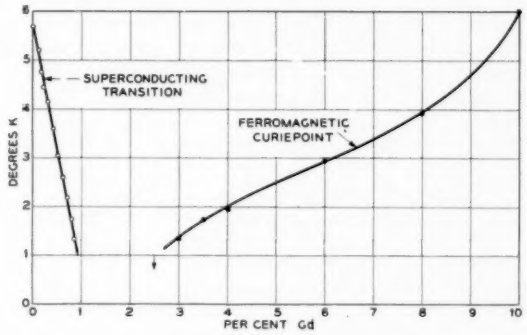

FIG. 3. Ferromagnetic and superconducting transition temperatures of solid solutions of gadolinium in lanthanum.

ture seems to be a strictly linear function of the amounts of dissolved gadolinium. 2.5 at.% or more of gadolinium in lanthanum causes this solid solution to become ferromagnetic above 1°K. The Curie points within this range are an approximately linear function of the percentage of gadolinium. This suggests the presence of a coupling which aligns the moments spontaneously in these materials and which is different from overlap exchange forces usually considered since the coupling extends over several lattice spacings and is proportional in magnitude to the amount of gadolinium added. By dissolving gadolinium in yttrium, a nonsuperconducting metal, only moderate paramagnetism was observed and solid solutions with even as much as 10 at.% gadolinium did not show any ferromagnetism. On the other hand, solid solutions of gadolinium in thorium, another superconductor, were again ferromagnetic.

These data suggest that an exchange over conduction electrons[2] leading to ferromagnetism is easy to bring about in an element which by itself

is a superconductor.[3]

In an earlier paper,[4] the question was raised whether an electronic configuration favorable to superconductivity might also be favorable to ferromagnetism. The experimental data by which we tried to illustrate this question started with the demonstration that while the germanides of most of the rare earths were either ferromagnetic or superconducting, all investigated silicides with the exception of $PrSi_2$, were neither ferromagnetic nor superconducting. We have now found that this dualism can also be illustrated by forming the AB_2 Laves phase compounds. If B is ruthenium then for A=Sc, Y, La, Ce we find superconductivity and for A=Pr···Gd···we have ferromagnets. In Table I the corresponding transition temperatures are listed.

Table I. Transition temperatures of ruthenides.

Compound	Behavior	T_c
$ScRu_2$	superconducting	1.67°
YRu_2	superconducting	1.52°
$LaRu_2$	superconducting	1.63°
$CeRu_2$	superconducting	4.9°
$PrRu_2$	ferromagnetic	40°
$NdRu_2$	ferromagnetic	35°
$ErRu_2$	ferromagnetic	13°

It appears rather remarkable that the compound $CeRu_2$, containing an element with a $4f$ electron, should be superconducting. However, the lattice constant reveals that in this compound at least 25% of the $4f$ electron has been squeezed into the outer shells.[5] In this way, the magnetic moment is reduced, and the number of valence electrons (hence the superconducting transition temperature) is raised considerably.

In summary, we wish to point out that there seems to be a correlation between the Curie points in certain alloys and compounds which include the rare earth elements and the superconducting transition temperatures of other closely related alloys and compounds.

Finally, we would like to thank P. W. Anderson and J. K. Galt for many helpful discussions.

[1] T. H. Van Vleck, Magnetic and Electric Susceptibilities (Oxford University Press, Oxford, 1932), pp. 243, 244.

[2] F. H. Spedding, Progress of Low-Temperature Physics (North-Holland Publishing Company, Amsterdam, C. Zener, Phys. Rev. 81, 440, (1951); S. V. Vonsorsky, J. Exptl. Thearet. Phys. U.S.S.R. Vol. II, p. 391; 981, (1946).

[3] Calculations by C. Herring, to be published in a subsequent paper, make this idea plausible.

[4] Matthias, Corenzwit, and Zachariasen, Phys. Rev. (in press).

[5] We want to thank Professor W. H. Zachariasen for his help and discussion with respect to these matters.

F AND V CENTERS THERMOLUMINESCENT RECOMBINATION*

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In the course of experiments on thermoluminescence of alkali halides, to be published in the near future,[1] the writers have found that current theories, all treating the process as a first-order one,[2] do not allow an easy interpretation of all the experimental features. It was decided, therefore, to investigate a particular case of thermoluminescence in which the nature of the electronic traps was known so that the task of setting up a model which could undergo quantitative examination would be greatly facilitated. X-rayed NaCl was chosen, in which F centers can give rise to thermoluminescence, as is well known.

Using an idea previously advanced by Johnson[3] and by Hill and Schwed,[4] a second-order process was considered and a phenomenological theory of F center thermoluminescence was carried out, starting from the equation

$$-dN_i/dt = p_i \exp(-E/kT)\, N_i(N_i - N_{i_0} + n_0). \quad (1)$$

This equation is based on the hypothesis that there are present in the crystal, together with the F centers, also some kinds of "luminophor" centers L_i. The electrons are supposed to be able to escape from the F-center ground level, with a thermal activation energy E, to reach some higher, excited level (perhaps in the conduction band) from which they can undergo a radiative transition to the L_i levels, with probability factors p_i. N_i means the luminophor center concentration, n the F-center concentration, the index 0 identifying the initial values at the beginning of the ith peak of the "glow curve" considered.

This way of reasoning can interpret satisfactorily many experimental features, and, moreover, it contains the following implications:

(a) The analysis of the glow curves must give a unique E value for all the peaks.[5]

(b) The light of each peak must have its own wavelength.

(c) The temperatures of the maxima of the glow curves are by no means characteristic of the trap depth, not even when the warming speed is held constant; on the contrary, they depend on the value of n_0 (in contradiction to the results of Randall and Wilkins' calculations[2]).

Experiments have been performed to test these points, with satisfactory results. It has been found: $E = 0.72$ ev for all the peaks,[6] and $p_1 = 5.10^{-10}$, $p_2 = 9.10^{-8}$, $p_3 = 2.10^{-5}$ cm^3 sec^{-1}.

In connection with point (b), we succeeded, through the use of interference filters, in proving that the wavelengths quoted by Hill and Schwed,[4] namely 3620, 4180/4320, and 5250 A, belong, respectively, to the first, second (doublet), and third peak.

The principal problem appeared then to be the identification of the L_i centers.

The following facts support the idea that L_i centers are V centers, introduced by irradiation at the same time as F centers.

First of all, it is well known that the relative amount of the different species of V centers is strongly influenced by the thermal history of the specimen, or by the irradiation temperature, and the same sensitivity is shown by the relative intensity of the thermoluminescence peaks. A much more important argument of a quantitative nature is that, if each L_i center is a V_i center, this means that when the radiative thermoluminescent transition takes place, the excited electron annihilates the V_i center, by recombination with the hole belonging to the V_i center. It can then easily be seen that the energy of the photon released in such a process is given by

$$(h\nu)_i = \epsilon - (h\nu)_{V_i} \qquad (2)$$

where ϵ means the energy gap between the excited level of the electron and the valence band, and $(h\nu)_{V_i}$ means the energy belonging to the V_i absorption band. Consequently, in Eq. (2), ϵ must have the same value for every thermoluminescence peak.

Since the $(h\nu)_i$ and $(h\nu)_{V_i}$ values can be found in the literature,[4,7] it has been possible to test our idea, and the results are given in Table I.

Table I shows that, with the exception of the first peak, our predictions are surprisingly well confirmed; and, in fact, the value found for ϵ is very near to the value of the energy gap of NaCl (about 9 ev).[8,9] The ϵ corresponding to the first peak deviates sensibly, but it is a pleasure to acknowledge here a kind private communication by Professor P. Pringsheim, who pointed out that the V_1 band energy value given in reference 7 was probably unreliable.

Table I. Values of ϵ calculated through $(h\nu)_i$ and $(h\nu)_{V_i}$.

i	1	2/3	4
$(h\nu)_i$	3.425	2.966/2.870	2.362
$(h\nu)_{V_i}$	3.594	5.570	5.904
ϵ	7.019	8.536/8.440	8.266

* The research reported in this document has been sponsored in part by the Air Force Office of Scientific Research of the Air Research and Development Command, U. S. Air Force, through its European Office.

[1] See also Bonfiglioli, Brovetto, and Cortese, Technical Notes No. 1b, 2b, 3b under Contract AF 61(514)-1333 of the ARDC European Office.

[2] F. Urbach, Akad. Wiss. Wien., Ber. (2A), 139, 363 (1930); J. T. Randall and M. H. F. Wilkins, Proc. Roy. Soc., (London) A184, 366 (1945).

[3] R. P. Johnson, J. Opt. Soc. Am. 29, 387 (1939).

[4] J. J. Hill and P. Schwed, J. Chem. Phys. 23, 652 (1955).

[5] As a matter of fact, the peaks found in our experiments are three, but the second is a doublet rather poorly resolved.

[6] Hill and Schwed's value is 1.28 ev.[4] This discrepancy has not been explained so far.

[7] Castler, Pringsheim, and Yuster, J. Chem. Phys. 18, 887 (1950).

[8] N. F. Mott and R. W. Gurney, Electronic Processes in Ionic Crystals (Oxford University Press, Oxford, 1940), p. 97.

[9] It is to be remarked that values for $(h\nu)_i$ and $(h\nu)_{V_i}$ in Table I are not obtained through measurements performed at the same temperature: $(h\nu)_i$ are thermoluminescent temperature values, whereas $(h\nu)_{V_i}$ are low temperature values. This point could be worth some discussion, but, nonetheless, the agreement obtained definitely proves our aim.

SHAPE AND POLARIZATION OF THE BETA SPECTRUM OF RADIUM E†

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The shape of the beta spectrum of RaE (Bi^{210}) has long been known to be nonallowed, even though it is a $\Delta I = 1$ (yes) transition in a heavy

element. An allowed shape usually occurs for such transitions because of the large Coulomb interaction which gives rise to leading terms in the spectrum shape which have the same energy dependence as for allowed spectra. In the case of RaE, the nonallowed shape has been explained as a result of the accidental cancellation between these usually dominant terms.[1]

The measurement of the longitudinal polarization of the beta particles affords a test of this rather surprising cancellation of the spectral shape parameters, since both experimental quantities depend on the same nuclear matrix elements. In the absence of cancellation, a polarization of nearly $-v/c$ would be expected. Recently, Geiger, Ewan, Graham, and MacKenzie[2] have measured an average polarization of the beta rays from a number of nuclei, and have found an anomalously small polarization in the case of RaE (50 to 80% of $-v/c$).

We have investigated the conditions the experimentally observed spectral shape imposes on the degree of polarization. In this analysis we have assumed the simplest formulation of the problem consistent with present experimental facts, namely: (1) the two-component theory of the neutrino, (2) invariance under time reversal, and (3) the V-A form of the beta-decay interaction.

The correction factor, $C(W)$, to the allowed spectral shape has been written in the form

$$C(W) = a_{-1} W^{-1} + a_0 + a_1 W + a_2 W^2. \qquad (1)$$

The degree of longitudinal polarization may be written[3]

$$P(W) = -\frac{v}{c}\left[1 - (a_{-1} W^{-1} + b_0)\frac{1}{C(W)}\right]. \qquad (2)$$

In these calculations we have included the effects of the finite nuclear size in the form given by Matumoto and Yamada.[3] The parameters a_i and b_0 appearing in (1) and (2) can be expressed in terms of five combinations of nuclear matrix elements, $M(j)$, instead of the usual three which appear in the point-nucleus approximation.

Using experimental values of the four parameters a_i, as determined from the measurement of the spectral shape, one can determine ranges of the five matrix element combinations, $M(j)$. These, in turn, imply a range of values of the degree of polarization. However, because the existing experimental determinations of the spectral shape are consistent with a great variety of sets of the coefficients a_i, and because of the preliminary nature of the existing polarization measurements, we report here only the results corresponding to a particular fit to the spectral shape - Plassmann and Langer's[54] "best fit" (their Fig. 3). Figure 1 presents the range of polarization consistent with this fit. The results in Fig. 1 are consistent with the value of the average polarization measured by Geiger, Ewan, Graham, and MacKenzie.[2]

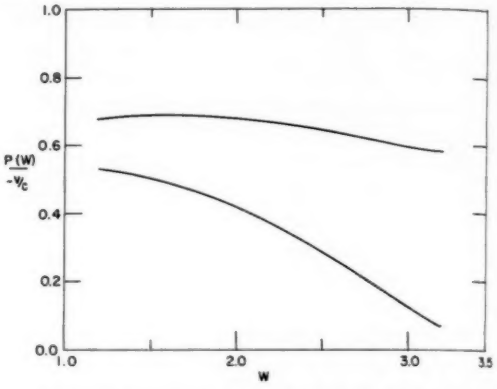

FIG. 1. Extreme limits of the longitudinal polarization of the electrons from the beta decay of RaE based on Plassmann and Langer's "best fit" to the spectral shape. W is the beta-particle total energy in units of $m_e c^2$. The range indicated results from the fact that the four shape parameters do not uniquely determine the five relevant matrix element combinations.

It should be emphasized that the degree of polarization is rather sensitive to the choice of the "fit" to the spectral shape. For example, for Plassmann and Langer's fit corresponding to their Fig. 5, the average polarization is raised to about 80% of $-v/c$. A measurement of the degree of polarization as a function of energy would provide a more sensitive test of the current interpretation of the beta decay of RaE. A paper is in preparation describing the details of the present calculations, including an interpretation of the matrix element combinations.

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*Guest scientist from Frankford Arsenal, Philadelphia, Pennsylvania.

[1] M. Yamada, Progr. Theoret. Phys. Japan 10, 252 (1953).

[2] Geiger, Ewan, Graham, and MacKenzie, Bull. Am. Phys. Soc. Ser. II, 3, 51 (1958).

[3] These results can be obtained, for example, from

the formulas of R. B. Curtis and R. R. Lewis [Phys. Rev. 107, 543 (1957)] with appropriate changes. The electron functions ($\kappa = \pm 1$) used are the solutions of the Dirac equation corresponding to a finite nuclear charge distribution, using the formalism of Z. Matumoto and M. Yamada [Progr. Theoret. Phys. Japan 19, 285 (1958)]. The less important electron functions ($\kappa = \pm 2$), which appear in the terms L_1 and L_1', have been taken as their "point nucleus" values. In particular, we assume $L_1 = c_0 + c_2 W^2$ and $L_1' = (p/W)L_1$, which are considered satisfactory approximations for the present purposes.

[4] E. Plassmann and L. Langer, Phys. Rev. 96, 1593 (1954).

K-MESON-NUCLEON INTERACTIONS AND THE RELATIVE Σ-Λ PARITY*

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In this note we wish to add to some recent theoretical considerations concerning the nature of the strong interactions of pions, nucleons, K mesons and hyperons.[1-6] The main conjecture that we want to make is the following: the K-Λ relative parity is odd relative to the nucleon and the K-Σ relative parity is even relative to the nucleon. In more specific terms, we conjecture that the interaction Hamiltonian has the form ($g_\Lambda \bar{N}\gamma_5 \Lambda K$ + Hermitian conjugate) for the Λ particle, nucleon, and K meson; and the form ($g_\Sigma \times \bar{N}\tau_a \Sigma_a K$ + Herm. conj.) for the Σ particle, nucleon, and K meson (the symbol for a particle denotes the operator that destroys it and the decomposition into charge states is to be understood). In the following we will make some arguments for our conjecture which will be based on perturbative estimates including certain higher-order effects.[2,7] These considerations are nonrelativistic in the sense that baryon-antibaryon pair states will not be considered; in other words, the γ_5 coupling will be reduced to the familiar P-wave vertex and the coupling with unity will be an S-wave vertex. The insights into low-energy K-meson phenomena so gained may compare favorably with the insights into low-energy pion phenomena gained with a similar approach in the early days of pion physics.

We consider first K^+ scattering by nucleons. In a recent calculation we estimated the scattering at low energies via a mechanism involving the exchange of a pion pair between the K meson and nucleon.[1,2] This mechanism gave rise to predominant S-wave scattering of K^+ by protons and neutrons equally and no charge exchange scattering. Call the amplitude for this scattering in either isotopic spin state, λ. This corresponds to an effective repulsive potential. Now let us introduce the scattering via the sequence $K+N \to K+K+\Sigma \to K+N$. This will be an attractive S-wave interaction. The amplitudes for the isotopic spin states zero and one will be $a_0 = g_\Sigma^2 \times C$, $a_1 = 3g_\Sigma^2 C$, where C is essentially the matrix element. Then the total amplitudes for the three reactions we are concerned with are as follows:

(1) $K^+ + p \to K^+ + p : \lambda - g_\Sigma^2 C$,

(2) $K^+ + n \to K^+ + n : \lambda - 2g_\Sigma^2 C$,

(3) $K^+ + n \to K^0 + p : g_\Sigma^2 C$.

These amplitudes are compounded from the total amplitudes for the two isotopic spin states, $A_0 = \lambda - 3g_\Sigma^2 C$ and $A_1 = \lambda - g_\Sigma^2 C$. We see that the S-wave scattering in the isotopic spin zero state is depressed. For example, if $\lambda \sim 3g^2{}_\Sigma C$, the charge exchange to non-charge exchange ratio would be ~ 0.2. If we now consider the scattering via the sequence $K+N \to K+K+\Lambda \to K+N$, we find that this will contribute to scattering in the P states. There will be an effective repulsion in the $P_{\frac{1}{2}}$ state of isotopic spin one, and an effective attraction in the $P_{\frac{1}{2}}$ state of isotopic spin zero (these signs are reversed for the $P_{\frac{3}{2}}$ states), the amplitudes being proportional to g_Λ^2. What is significant here is that of the two "small" cross sections at low energies, corresponding to reactions (2) and (3) above, the P-wave scattering through the intermediate Λ state will contribute in lowest order only to the charge exchange process. This could account for the experimentally observed striking rise in the charge exchange to non-charge exchange ratio at the same time as P-wave scattering appears to be making its appearance in the isotopic spin zero state.[8,9] That the P-wave scattering should be felt first in the isotopic spin zero state is, of course, a natural consequence of the depressed S-wave scattering in this state relative to that in the isotopic spin one state, but we should expect some S-P interference in K^+-p scattering as well, at the higher energies now under study.

We now want to remark upon the surprising infrequency of Λ particles produced in low-energy K^--p interactions relative to any of the Σ particles produced by this reaction.[10] If one accepts the conjecture of odd relative Σ-Λ parity, one readily sees that very low-energy S-wave K^--p

interactions will produce the Σ-π system (in an isotropic $P_{\frac{1}{2}}$ state), whereas the Λ-π system (in a $P_{\frac{1}{2}}$ or $P_{\frac{3}{2}}$ state) will result from the interaction of P-wave $\bar{K}$. Thus, we would expect the Λ production on free protons to become more frequent as the K^- energy is raised, and also to be more frequent in interactions with protons in heavy nuclei, as appears to be the case experimentally.[11]

Finally, we remark that the absence of a strong polarization of the Σ particles produced in $\pi^- - p$ collisions at ~1 Bev, as contrasted with the strong polarization of the Λ particles produced in these collisions,[12] may have a simple origin in the absence and presence, respectively, of a dominant spin dependence in the $N\Sigma K$ and $N\Lambda K$ vertices.

If the conjecture presented here is correct, there would certainly be a difference in the dynamics of the various meson-baryon and baryon-baryon interactions, even if the neglect of certain phenomena depending on baryon mass differences happens to be a reasonable approximation. There might still be an over-all strength of "bare" coupling approximately common to what we call the strong interactions.

I wish to thank Dr. F. Gursey for a vital discussion upon the possible theoretical formulations of these ideas, other aspects of which Dr. Gursey will discuss in a forthcoming paper. I wish to express sincere appreciation to Professor Abraham Pais for a number of very stimulating discussions and for his continuing advice and encouragement.

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*Work performed under the auspices of the U. S. Atomic Energy Commission.

†A National Science Foundation Postdoctoral Fellow.

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POSSIBLE CONNECTION BETWEEN STRANGENESS AND PARITY*

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In this note it will be shown that a connection between strangeness and parity can be conjectured in such a way that predictions become possible on the relative parities of strongly interacting particles. Some connection between the two conservation laws has been suspected and discussed by several authors.[1-3] Here we wish to point out that the idea due to Yang and Tiomno,[4] of considering four kinds of spinor fields with respect to their behaviour under P (parity), seems to provide the generalization of the parity operation ($P^2 = \pm 1$) needed for establishing a one-to-one correspondence between the isotopic parity and the space parity of charge multiplets. This leads one to assign opposite parities to Λ and $\bar{\Sigma}$ as well as to N and Ξ for baryons, and to K and $\bar{K}$ for mesons.

Under P, the 4-spinor fields of types A, B, C, and D transform respectively according to the laws

$$\vec{r} \to -\vec{r};\ \psi_A \to \gamma_4 \psi_A,\ \psi_B \to -\gamma_4 \psi_B,$$
$$\psi_C \to i\gamma_4 \psi_C,\ \psi_D \to -i\gamma_4 \psi_D. \quad (1)$$

It follows that by combining such spinors, one can also construct four types of Lorentz invariant functions ϕ, which, under P transform according to the laws

$$\vec{r} \to -\vec{r};\ \phi_A \to \phi_A,\ \phi_B \to -\phi_B,$$
$$\phi_C \to i\phi_C,\ \phi_D \to -i\phi_D. \quad (2)$$

We note immediately that, although a real scalar ϕ_A and a real pseudoscalar ϕ_B can be constructed, the fields ϕ_C and ϕ_D are necessarily complex. Hence, a neutral boson field that is identical with its charge conjugate (up to a sign difference) will belong to the types A, B, whereas a neutral boson field different from its charge conjugate will be described by the types C, D. Using charge independence, we assume that fields describing members of the same charge

multiplet belong to the same type. We further postulate that each of the eight classes under consideration is represented by a different charge multiplet that we must identify with physical particles. We distinguish between the types C, D (fields with hypercharge) and A, B (fields with no hypercharge) by the fact that $P^2 = -1$ and $P^2 = +1$, respectively, for the two groups. Each group contains two classes of opposite parities with opposite values of the hypercharge. Because of the behaviour of meson fields under charge conjugation and the pseudoscalar nature of the pion $\vec{\pi}$, we make the following assignments:

$$\phi_B = \vec{\pi}, \quad \phi_C = K, \quad \phi_D = \bar{K}, \tag{3}$$

where K has components K^+, K^0 and $\bar{K} = i\tau_2 K^*$. ϕ_A corresponds to an unobserved particle, namely the neutral isosinglet scalar meson π_0' discussed in the literature. It also follows from (3) that the parity of K cannot be compared directly to the parity of $\vec{\pi}$, in agreement with Pais's remarks[1] about the intrinsic parities of strange particles. If $K = K_1^0 + iK_2^0$, then, under P, we have

$$\vec{r} \to -\vec{r}; \quad K_1^0 \to -K_2^0, \quad K_2^0 \to K_1^0. \tag{4}$$

We now identify the baryon fields N, Ξ, Λ, and $\vec{\Sigma}$ with the classes defined by (1). Since K and N must have the same hypercharge to give the observed strangeness selection rules, we identify ψ_C as N, so that the other doublet Ξ must belong to the class D. Classes A and B have to be associated with the remaining charge multiplets, namely Λ and $\vec{\Sigma}$. Either identification ($\Lambda = \psi_A$, $\vec{\Sigma} = \psi_B$) or ($\Lambda = \psi_B$, $\vec{\Sigma} = \psi_A$) is possible. Tentatively choosing the second alternative, we have

$$\psi_A = \Sigma, \quad \psi_B = \Lambda, \quad \psi_C = N, \quad \psi_D = \Xi. \tag{5}$$

Here we stress the fact that N can belong to the Class C without necessarily having a Majorana structure, the function obtained by combining a spinor of type A and a boson of type C being such an example. More generally it is easy to construct isotopic multiplets with the assignments (3) and (5) if one expresses all these wave functions as combinations of four basic spinor fields with suitable group properties, also using the fact that γ_5 switches the parity of a 4-spinor. Apart from the factor γ_4 and a minus sign for $\vec{\pi}$ and Λ, the parity factors in Eqs. (1) and (2) correspond to the Racah[5] parity assignments in isotopic spin space. The parity operator for baryons is

$$P\psi(t, \vec{r}) = P_R \gamma_4 \psi(t, -\vec{r}), \tag{6}$$

where P_R is our modified isotopic spin parity. P_R, unlike the Racah parity, changes sign on multiplying the baryon wave function by γ_5, since now P_R is not defined independently of the space parity. For example $\vec{\Sigma}$ and $\gamma_5\Lambda$ have the same P_R. Now, if one writes an interaction Lagrangian of the type (direct Yukawa couplings) considered by D'Espagnat and Prentki,[6] and one requires charge conservation and invariance under P, then one is led not only to the strangeness selection laws, but to a more restricted expression where the place of γ_5 in the interaction terms becomes uniquely fixed. Assuming, for simplicity, a common coupling constant g_π for pion interactions and g_K for K-interactions, and also invariance under T, we obtain for example

$$\begin{aligned} L_{\text{int}} = {} & g_\pi (i\bar{N}\gamma_5 \vec{\tau} N \cdot \vec{\pi} + i\bar{\Xi}\gamma_5 \vec{\tau} \Xi \cdot \vec{\pi} + \bar{\Lambda}\vec{\Sigma} \cdot \vec{\pi} \\ & + i\vec{\Sigma} \times \gamma_5 \vec{\Sigma} \cdot \vec{\pi}) + g_K (i\bar{N}K\gamma_5\Lambda + \bar{N}\vec{\tau}K \cdot \vec{\Sigma} \\ & + i\bar{\Xi}\gamma_5 \bar{K}\Lambda + \bar{\Xi}\vec{\tau}\bar{K} \cdot \vec{\Sigma}) + \text{Herm. conj.} \end{aligned} \tag{7}$$

We note that invariance under P alone would still lead to charge conservation for the above Lagrangian but would also permit the charge violating term $\bar{N}\vec{\tau}\Xi \cdot \vec{\pi}$. Dropping the Yukawa restriction, there is also the possibility of adding a term[7] of the form $K^\dagger K \vec{\pi} \cdot \vec{\pi}$. On the other hand, $(\bar{K})^\dagger \vec{\tau} K \cdot \vec{\pi}$ is ruled out by charge conservation.

It is remarkable that Katsumori[8] has obtained the correct splitting of baryon masses to the first order by assuming that Ξ and N as well as Λ and $\vec{\Sigma}$ have opposite parities. However, because of our unconventional definition of P, our (Ξ, K) terms in (7) differ from Katsumori's Lagrangian as to the place of γ_5. Hence his arguments for the ($N\Xi$) mass splitting will not apply in this case and a different coupling constant for the (ΞK) interaction may be necessary. Finally, as stressed by Barshay,[9] indirect support for Λ and $\vec{\Sigma}$ having opposite parities is provided by the analysis of hyperon production and K-N scattering experiments. Other experiments under way will probably give more direct evidence on the relative parities of baryons and test the hypothesis of an intimate link between parity conservation and strangeness conservation.

I would like to express my gratitude to Dr.

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INVESTIGATION OF TIME-REVERSAL INVARIANCE IN THE BETA DECAY OF THE NEUTRON

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It has been pointed out by Jackson *et al.*[1] that a violation of time-reversal invariance in the beta decay of the neutron would give rise to a term of the form $D\vec{J}\cdot(\vec{p}_e\times\vec{p}_\nu)/JE_eE_\nu$ in the angular distribution of the decay products. In this expression J is the spin of the neutron; p_e, E_e, and p_ν, E_ν are the momenta and energies of the electron and neutrino, respectively. The coefficient D of this term is related to the beta-decay coupling constants by the expression

$$D=\frac{2\,\mathrm{Im}(C_SC_T^*-C_VC_A^*+C_S'C_T'^*-C_V'C_A'^*)}{|C_S|^2+|C_V|^2+|C_S'|^2+|C_V'|^2+3(|C_T|^2+|C_A|^2+|C_T'|^2+|C_A'|^2)}$$

We are measuring the magnitude of this coefficient D using a beam of neutrons from the NRX reactor as a source.

Figure 1 shows a plan view of the apparatus mounted adjacent to the shield of the reactor.

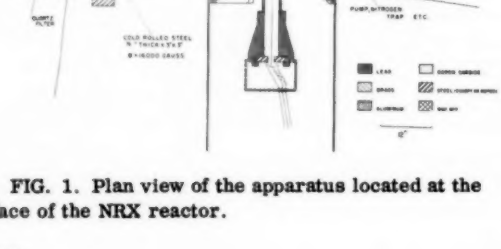

FIG. 1. Plan view of the apparatus located at the face of the NRX reactor.

The collimated beam passes through a 3.5-inch thick filter of crystalline quartz where the unwanted epithermal neutrons are preferentially scattered out prior to polarization of the thermal neutrons. The quartz, held at room temperature, causes a reduction in the epithermal neutrons of a factor of 3.2 and a reduction in the thermal neutrons of a factor of 1.9. The filtered beam then passes through a block of cold rolled steel, one-half inch thick, magnetized in a vertical direction to a flux density greater than 16 000 gauss. The preferential scattering by the iron of those neutrons whose spin direction is parallel to the direction of the magnetic field results in the transmitted beam being partially polarized. After further collimation and in the presence of a magnetic guiding field of about 150 gauss, the polarized neutrons enter a vacuum tank in which decays are observed by the simultaneous detection of decay

electrons and recoil protons. After traversal of the vacuum tank the beam is monitored by a fission chamber located in the beam catcher.

The region of the neutron beam from which decay events are detected is surrounded by an electrostatic shield held at +7 kv. Protons from decay events occurring in a portion of the beam about one inch in length which recoil into a 10-degree cone about a line perpendicular to both the beam direction and the direction of polarization pass through a wire grid and, their direction established, are first accelerated through 1000 volts and then further accelerated and focussed by a 6000-volt cylindrical electrostatic lens onto the first dynode of an electron multiplier. The electrons are detected by two scintillation counters using 5-in. diameter by $\frac{1}{8}$-in. thick plastic phosphors. The resolution of the scintillation counters is about 30% and their outputs are passed through pulse-height analyzers to select those pulses corresponding to electrons of energy between 350 and 550 kev. The electron counters are located at 160° from the proton direction to maximize the product $(\vec{p}_e \times \vec{p}_\nu)$.

Coincidences between the proton counter and the two electron counters separately are recorded, and the asymmetry between the two coincidence rates is determined for both polarized and unpolarized neutrons. The latter condition is achieved by switching off the polarizing magnetic field with the guide magnetic field still on. Under this condition the magnetic field over the region of the electron trajectories resulting from the stray field of the guide system is always present. This stray field amounts to about 4 gauss and is unaffected by the polarizing field. The difference between the asymmetries with polarized neutrons and with unpolarized neutrons is used to obtain the time reversal coefficient D.

Most of the data has been accumulated using a one-half inch iron block as polarizer and with an arrangement which records coincidences for one hour with polarized neutrons followed by one hour with unpolarized neutrons in an automatic sequence. Under these conditions about 18 true events and 35 random events are recorded in each of the two coincidence channels for unpolarized neutrons and a similar number for polarized neutrons in 20 hours of operation. Data have also been taken without the quartz filter and also with a one inch block of iron as a polarizer. Checks have also been made on the asymmetry with no quartz or iron in the beam to provide a high-intensity source of unpolarized neutrons.

The polarization was measured without disturbing the magnetic fields in the region of the neutron beam by diffracting 2.8-A neutrons from the (100) planes of a cooled Haematite antiferromagnetic crystal located at the source volume of the neutron beam and then diffracting these neutrons again from the (111) planes of a magnetized face-centered cubic crystal of cobalt iron alloy. This showed that the direction of polarization was within five degrees of the vertical and with the aid of the method of Stanford *et al.*[2] indicated that the average polarization of the beam was about 40% with the one inch iron block as polarizer. By comparison of the single transmission effects the polarization was estimated to be 27% with the half-inch block and the quartz filter.

Combining the measurements under the various experimental conditions yields a present value of $D = -0.02 \pm 0.28$ where the error is primarily statistical but includes an allowance for the uncertainty in polarization and geometric corrections. This value has to be compared with expected values of $D = 0$ for time-reversal invariance not being violated and $|D| = 0.5$ for full violation with an equal mixture of axial vector and vector interactions each with $C_i' = C_i$. The present result is thus consistent with time-reversal invariance and suggests that full violation does not occur. The experiment is continuing but due to the current interest in this topic it is through that this preliminary result may be useful.

* Supported in part by the U.S. Atomic Energy Commission and the National Security Agency.

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EXPERIMENTAL LIMIT OF THE NEUTRINO REST MASS*

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As stated recently,[1] an empirical determination of the neutrino rest mass can be obtained from the H^3 - He^3 mass difference and the H^3 β

end-point energy. Sakurai[2] has noted that this is of course valid only if one has the true end-point energy but not for the extrapolated value. The spread in β end-point data reported in the literature, however, is large, i.e., 1.3 kev, when compared with the magnitude of extrapolation from the last experimental points in any one of the β spectra. Consequently the difference between the experimental lower limit for the true β end point and extrapolated end point is comparable with or less than experimental error and at present trivial. If one takes the experimental lower limit of the true end point, i.e., the last experimental point in the β spectrum, and subtracts this from the H^3 - He^3 mass difference, the result is an empirical upper limit for the neutrino rest mass. This upper limit is independent of any theory of β decay and requires only the assumption of conservation of mass and energy. This was the basis of the estimate of the order of 1 kev for the upper limit of the neutrino rest mass by Friedman and Smith.

The considerations of Sakurai will become relevant only when the β end-point (true or extrapolated) can be measured with much higher accuracy than it is now known.

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EXPERIMENTAL EVIDENCE FOR THE INFLUENCE OF ATOMIC BINDING ON THE DECAY RATE OF NEGATIVE MUONS*

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The entire body of observations on the rate of disappearance (Λ_{dis}) of negative muons in matter,[1] in particular the striking Z dependence of this rate, can be described by attributing the disappearance to a competition between nuclear capture (at a rate Λ_{cap}) and decay (at a rate Λ_{dec}), both taking place from the mesic K orbit. Λ_{cap} depends primarily on the probability of finding the muon inside the nucleus (Z, A) and thus strongly on Z (Wheeler's Z^4 law); on the other hand, one assumes customarily that Λ_{dec} is uninfluenced by the fact that the decay takes place from a bound orbit, i.e., one puts $\Lambda_{dec}(Z) = \Lambda_{dec}(0) \equiv \Lambda_{dec}(\mu^+)$. This assumption is in particular made when capture rates are calculated from the experimentally accessible disappearance rates.[2]

It is, however, clear that the ratio $R \equiv \Lambda_{dec}(Z)/\Lambda_{dec}(0)$ should differ from unity for several reasons: (a) reduction of the number of final states through the decrease in available energy; (b) Doppler effect and time dilatation caused by the motion in the K orbit; (c) distortion of the outgoing electron wave by the nuclear electric potential. Effects (a) and (b), which jointly lead to $R<1$, have been considered by various authors[3] for the case of point-like nuclei, while effect (c), which works in the opposite direction, has so far been consistently neglected.

In view of this situation and in order to enable one to make precise determinations of Λ_{cap} from Λ_{dis}, we are now measuring R directly for a variety of elements. Some preliminary results will be given here.

Our measurements are based on the following: Two quantities are, in principle, directly measurable: (1) $\Lambda_{dis}(Z, A)$, and (2) $N_e(Z)$, the number of electrons per N_μ muons stopping in a target (Z, A). As $N_e = N_\mu[\Lambda_{dec}(Z)/\Lambda_{dis}(Z, A)]$, an absolute measurement of these quantities would yield R directly. Such a measurement is however difficult as far as N_e is concerned. On the other hand, it is sufficient to count the N_e's from pairs of targets (Z_1, Z_2) with the same relative efficiency to determine the Z dependence of R. With the assumption that, say $\Lambda_{dec}(Z_1) = \Lambda_{dec}(0)$, R itself may be determined.

"Sandwich" targets were built up from alternating layers of elements Z_1 and Z_2. Each comprised about 5 sheets of either material and had a total thickness (6 - 8 g/cm²) sufficient to stop most of the local high-purity μ^- beam.[4] The ratio $S \equiv N_\mu(Z_1)/N_\mu(Z_2)$ was determined directly with an anticoincidence range telescope that included, just behind the target, a thin (1/16 in.) scintillator. S could be obtained to an accuracy of ± 2%. The ratio $N_e(Z_1)/N_e(Z_2)$ was measured by observing the composite time distribution of decay electrons from the sandwich by means of the time-converter arrangement which had also served to measure $\Lambda_{dis}(Z)$ for pure targets.[2] A coincidence telescope consisting of three scintillators separated by two 1.8-g/cm² Al absorb-

ers was used now to detect the electrons. Such a telescope has a negligibly small detection efficiency for gamma rays of less than 10-Mev energy. This was essential for a correct determination of R, as a large number of gamma rays, N_γ, appear to be produced per muon capture, presumably from the de-excitation of capture products. N_γ may, in particular in the case of heavy nuclei, exceed N_e by a factor 100.

Table I summarizes the results obtained so far. The errors quoted for R make some allowance for systematic uncertainties. These stem mainly from two causes: (1) differences in the shape of the decay electron spectra[5] from Z_1 and Z_2, (2) differences in the value of the asymmetry parameter a which governs the angular distribution of the electrons. We have some evidence, obtained by varying target thickness, that (1) is probably unimportant. While no good determinations of a are available for the targets and the beam used here, it may be anticipated that a will be nonzero only for spinless nuclei and in their case will not exceed 0.06, the local value for $Z = 6$ in good geometry. Thus (2) should be unimportant also, in particular when the strong multiple scattering in the thick targets used is taken into account.

It appears from our results that R is, with the exception of Fe ($Z = 26$), less than unity. For this element, effects (a) and (b) would give $R \cong 0.80$. It may be inferred that effect (c) dominates near this Z value,[6] but is overshadowed by the other two effects as Z increases. While the trend of $R(Z)$ found here is not unreasonable, comparison with theory must await the completion of computations using exact muon and electron wave functions.[7]

A practical conclusion from our measurements is that the corrections to Λ_{cap} do not exceed a few percent, when Λ_{dis} is the experimentally measured quantity. The determination of Λ_{cap} from N_e, such as was used by Lederman and Weinrich,[8] involves however corrections directly proportional to R, as was recognized by these workers themselves.

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Table I. Determination of R for various elements.[a]

Sandw.	Z_1/Z_2	S	$\frac{N_e(2)\Lambda_{dis}(2)}{N_e(1)\Lambda_{dis}(1)}$	Lifetimes used ($\times 10^{-9}$ sec)	R[c]
C-Al[b]	6/13	0.543 ± 0.015	1.82 ± 0.05	$\frac{2010 \pm 20}{880 \pm 20}$	0.99 ± 0.04[d]
C-Fe	6/26	0.218 ± 0.005	5.57 ± 0.14	$\frac{2010 \pm 20}{202 \pm 7}$	1.21 ± 0.05
C-Zn	6/30	0.136 ± 0.004	7.09 ± 0.20	$\frac{2010 \pm 20}{161 \pm 6}$	0.96 ± 0.05
Al-Mo	13/42	0.469 ± 0.010	1.98 ± 0.06	$\frac{880 \pm 20}{105 \pm 4}$	0.93 ± 0.05[e]
C-Cd	6/48	0.121 ± 0.003	7.42 ± 0.30	$\frac{2010 \pm 20}{96 \pm 4}$	0.90 ± 0.05
C-W	6/74	0.175 ± 0.008	3.03 ± 0.20	$\frac{2010 \pm 20}{81 \pm 4}$	0.53 ± 0.05

[a] We assume throughout that the process $N + \mu^- = N + e^-$ does not occur; see S. Lokanathan and J. Steinberger, Phys. Rev. 100, 1490 (1955).

[b] $Z = 6$ in the form of CH_2 sheets; assumed to behave like elemental C.

[c] $R(6)$ assumed to be unity.

[d] Done to check consistency of method, as $R(13) = 1$ is expected.

[e] $R(13)$ assumed to be unity.

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[4]N. Campbell and R. A. Swanson (to be published).

[5]Reference 3, in particular Muto *et al.*

[6]J. C. Sens *et al.*, Bull. Am. Phys. Soc. Ser. II, 3, 198 (1958); recent preliminary results obtained at Liverpool also seem to indicate that $R>1$ near $Z = 26$ [H. Muirhead (private communication)].

[7]Numerical calculations yielding such functions are now in progress at Los Alamos Scientific Laboratory.

[8]L. Lederman and M. Weinrich, Proceedings of the CERN Symposium on High-Energy Accelerators and Pion Physics, Geneva, 1956 (European Organization of Nuclear Research, Geneva, 1956), Vol. 2, p. 427.

COSMIC-RAY INCREASES PRODUCED BY SMALL SOLAR FLARES*

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This note reports the first observations of two unusually short-lived increases in the low-energy portion of cosmic rays. They appear to be produced by small solar flares. Some indications of effects of small solar flares have been previously reported,[1] but have not been widely accepted, possibly because the connection was only statistical and an isolated increase could not be discerned.

We therefore designed an experiment whose purpose it was to detect short-lived or small increases in the low-energy component of total cosmic-ray intensity, observe their structure and study their correlation in time with small solar flares. In order to improve the chances of detecting such events we constructed a detector with very large sensitive area so that the counting rate was high enough to be recorded directly through a ratemeter. Furthermore, since it seemed likely that the low-energy portion of cosmic rays would be most affected, we operated this detector at high latitudes (55° geomagnetic) and at high altitudes (up to 45 000 feet) in an aircraft operated by the Rome, New York, Air Force Base.

The detector was a cosmic-ray counter telescope 1 square meter in area. It consisted of two trays of geiger tubes spaced 1 inch apart. The trays were operated in coincidence. No absorber was used.

Flights were made at various altitudes on August 8th and 9th, 1957. No rapid fluctuations in the total cosmic-ray intensity were observed during the three hours the aircraft was at altitude on August 8th. On the 9th of August, however, two unusual events were observed.

The first event (see Fig. 1) occurred while the aircraft was flying straight and level at an altitude of 25 000 feet. It started at 1349 UT and lasted for nearly three minutes. During this time the intensity increased from its normal value by approximately 30%. The second event started at 1433 UT while the aircraft was still at an altitude of 25 000 feet. It lasted for almost

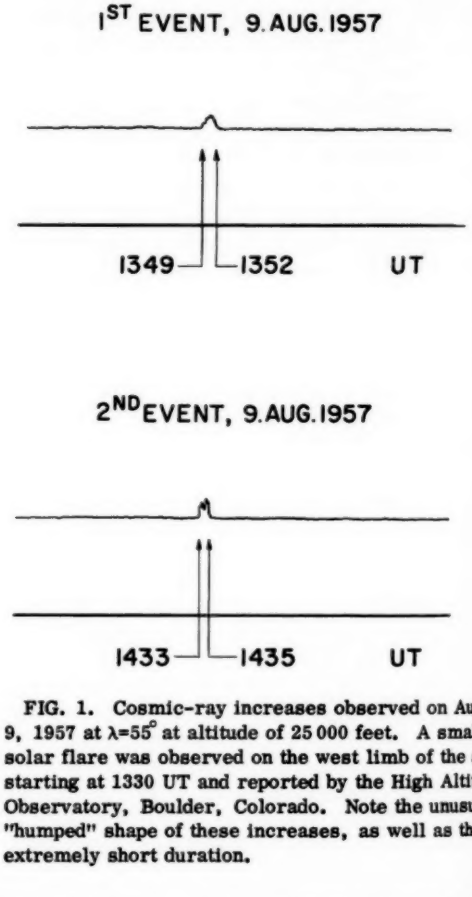

FIG. 1. Cosmic-ray increases observed on August 9, 1957 at $\lambda=55°$ at altitude of 25 000 feet. A small solar flare was observed on the west limb of the sun starting at 1330 UT and reported by the High Altitude Observatory, Boulder, Colorado. Note the unusual "humped" shape of these increases, as well as their extremely short duration.

two minutes. The intensity increase was again approximately 30%. The average coincidence rate was about 1000 counts per second.

Reports of solar and other geophysical activity were obtained for the two days the flights were made. No solar flares were observed on the 8th of August during the time measurements were made. On August 9th, a solar flare of importance 1 was observed on the west limb of the sun at 1330 UT; the activity lasted for about one hour (High Altitude Observatory, Boulder). Thus the onset of the flare preceded the first cosmic-ray event by about 20 minutes. Other geophysical events were observed during the same period. A radio noise storm on 200 Mc was in progress when the flare started and during the first cosmic-ray increase (reported by National Bureau of Standards, Boulder). A sudden ionospheric disturbance with shortwave fadeout (importance +3) followed the flare (reported by Central Radio Propagation Laboratory of the National Bureau of Standards, Boulder). A distinctive magnetic event was observed (at Fredericksburg, Virginia, Station) just at the time of the first cosmic-ray increase.

Several conclusions can be drawn from the results of these flights. The results offer the first direct experimental evidence for a time correlation between a small solar flare and cosmic-ray intensity increases of short duration. The increases followed the appearance of the flare by 20 minutes. The same time interval has been found between large solar flares and large cosmic-ray increases. But the increases reported here were in the form of a hump instead of a sharp rise and exponential decrease to normal.[2] The absence of similar increases at other cosmic-ray stations is indicative of the association of low-energy cosmic-ray particles with small solar flares. Evidently most of the particles causing the intensity increases have low energies and could not have reached sea level.

It seems probable therefore that many, perhaps most, solar flares can accelerate particles to energies which may however be quite low.[3] Provisionally we would assume that the low-energy cutoff ("knee") observed for the general cosmic radiation does not exist for the solar-flare cosmic rays. No statement can as yet be made about their charge spectrum.

The present experiment clearly shows the need for operating high counting rate cosmic-ray detectors in conjunction with short resolving time count ratemeters at high altitudes and at auroral latitudes to detect low-energy cosmic-ray particles associated with small solar flares. Failure of other cosmic-ray experiments to meet these two conditions may be the reason why such events have not been detected in the past.

Equipments similar to the one flown at high altitudes are at present in operation at mountain altitudes at Climax, Colorado (altitude 3400 meters), and at Banff, Canada (altitude 2283 meters) as part of the International Geophysical Year effort. Balloon flights are planned for the near future.

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*Supported by the Air Force Office of Scientific Research. We are particularly indebted to Mr. Ray Heer for his continued interest in this program and for making the difficult arrangements for the airplane flights.

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FREQUENCY OF CESIUM IN TERMS OF EPHEMERIS TIME

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(Received July 7, 1958)

The National Physical Laboratory, Teddington, and the U. S. Naval Observatory, Washington, have been cooperating in a joint program since June 1955 to determine ν_E, the frequency of cesium in terms of the second of Ephemeris Time.[1] In 1955 the International Astronomical Union recommended that the second of Ephemeris Time be adopted as the fundamental unit of time, and in 1956 the International Committee of Weights and Measures redefined the second so as to make it identical with the second of Ephemeris Time (E.T.), which is considered to be a constant unit of time. The second of Universal Time is thus no longer the fundamental unit of time.

Ephemeris Time is defined by the orbital mo-

Table I. Results for ν_E obtained from four different sets of data.

		Means	ΔH (sec)	$\nu_E - \bar{\nu}_U$	ν_E	$\Delta T''$ (sec/yr^2)
1.	ΔT_0,	1954.25-1958.25	+1.146	-121	9 192 631 761	+ 0.17
2.	ΔT_0,	1955.25-1958.25	1.085	-115	767	+ 0.10
3.	ΔT_C,	1954.25-1958.25	1.035	-110	772	+ 0.12
4.	ΔT_C,	1955.25-1958.25	0.966	-102	780	+ 0.17

tion of the earth about the sun, but is obtained in practice from the orbital motion of the moon about the earth. The dual-rate moon position camera has been used at the U.S. Naval Observatory since June, 1952, to determine Ephemeris Time.[2] Photographs of the moon and surrounding stars are taken to determine the position of the moon at a known Universal Time. The Improved Lunar Ephemeris 1952-1959 tabulates the position of the moon as a function of E. T. The Lunar Ephemeris is entered with the observed position of the moon and E. T. is taken out. There is thus obtained the quantity ΔT = E.T.-U.T.

Semiannual means have been determined for the epochs 1952.75 to 1958.25, in two forms. ΔT_0 is the mean obtained initially and ΔT_C is the mean obtained by correcting for terms which depend upon the mean anomaly of the moon.

In a previous note[1] we have described an intermediate step, namely, the determination of ν_U, the frequency of cesium in terms of the second of UT2. In this note we describe how the observations of the moon are used to convert from ν_U to ν_E.

It may be shown that

$$\nu_E = \bar{\nu}_U - \nu_E (\Delta H/H),$$

where ν_U is the mean value of ν_U in an interval of time H, and ΔH is the total change in ΔT during the interval.

The interval of comparison used is 1955.50 to 1958.25. For this interval $H = 8.68 \times 10^7$ sec, and $\bar{\nu}_U$ = 9 192 631 822 cps of UT2. ΔH was obtained by passing a parabola through the means, by least squares, and evaluating ΔT for the ends of the interval. Four solutions were made in order to determine the effect of using different data. The results are shown in Table I.

The last column gives the deceleration in the rotation of the earth as determined with the moon camera. The deceleration previously determined with the cesium standard was 50 parts in 10^{10} per year, which is equivalent to 0.16 sec/yr^2. The agreement between the moon camera and cesium is satisfactory.

The value of ν_E adopted is the mean of the four solutions, that is, 9 192 631 770 cps.

With this value of ν_E we may obtain ΔA = A.T. -U.T., where A.T. denotes atomic time, by integrating the values of ν_U previously obtained and by assigning an initial value of ΔA. The value adopted is ΔA=30.580 sec at 1957.0. Figure 1 shows ΔA, ΔT_0, ΔT_C, and ΔT_{C3}, the parabola obtained from solution No. 3. Its equation is

$$\Delta T_{C3} = [30.853 - 0.469\,(t - 1950.0) + 0.0615\,(t - 1950.0)^2]\ \text{sec.}$$

It is believed that the dispersion in the values of ΔT will be reduced by the application, in the future, of corrections for the figure of the moon which are now being completed by Dr. C. B. Watts of the Naval Observatory.

The probable error of ν_E is estimated to be ±10 cps from internal considerations. The result, however, may be affected by possible systematic errors, especially in the determination of ΔH. There is also the possibility that a systematic error may be present in the chain of

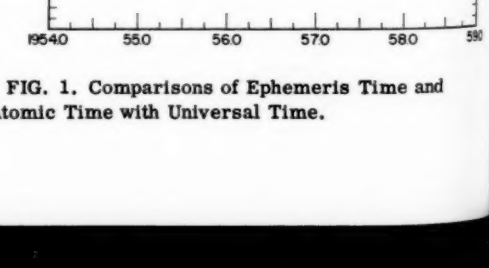

FIG. 1. Comparisons of Ephemeris Time and Atomic Time with Universal Time.

comparison between the cesium beam at Teddington and the moon camera at Washington. From an analysis of the various factors involved we have adopted a probable error of ±20 cps.

We find, thus, the transition frequency of cesium (4,0) — (3,0) at zero magnetic field is

$$\nu_E = 9\,192\,631\,770 \pm 20 \text{ cycles per second (of E.T.) at 1957.0.}$$

The mean epoch is specified because there is a possibility that the gravitational and atomic time scales may not be the same, and may change secularly. Future determinations of ν_E will decide this question.

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FURTHER OBSERVATIONS ON THE NATURE OF THE CURRENT REDUCTION IN THE PRIMARY COSMIC-RAY INTENSITY*

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It has been established recently that a sharp cutoff mechanism, operative now but not in 1949–50, has reduced the primary cosmic-ray intensity.[1] Although the experimental results which led to this conclusion were entirely consistent with the possibility that the 1958 and 1950 magnetic ridigity spectra differ only with respect to the lower limit, it was not possible on the basis of the available data to determine whether the form of the spectrum, *i.e.*, the exponent γ in the power law representation $J(>pc/Ze)=K(pc/Ze)^{-\gamma}$, had also changed.

This is a preliminary report upon the results of experiments with balloon-borne pulsed ionization chambers, which not only relate to this matter, but which also provide information regarding the flux of heavy nuclei.

An appropriately biased pulsed ionization chamber instrument responds principally to (1) nuclear disintegrations produced in the walls by fast protons and neutrons, and (2) heavy nuclei of atomic number $Z>8$. A detailed description of the experimental arrangement utilized in the present investigations, and a discussion of the interpretation of the measurements, has been presented previously.[2]

At balloon altitudes, events of Type (1) are produced predominantly by high-energy primaries and their progeny. Low-energy protons (*i.e.*, close to the geomagnetic cutoff at $\lambda \approx 52°$N) yield a negligible contribution to the disintegration rate. This was demonstrated by the fact that, in 1950, the burst rate did not change between $\lambda = 52°$ and $\lambda = 69°$ (statistical uncertainty ±3%, atmospheric depths 7 -70 g/cm²) although the vertical primary intensity increased by 45%.[2]

However, it was not possible to deduce any quantitative information regarding the low-energy portion of the heavy primary spectrum because most particles of atomic number $Z>8$ in the band of energies admitted between $\lambda = 52°$ and $\lambda = 69°$ are absorbed as a result of energy loss by ionization in the residual atmosphere about the apparatus. This sharp atmospheric cutoff effect, somewhat surprising at that time, is now understandable in terms of the discrepancy between geomagnetic coordinates and directly measured minimum cutoff energies.[3] However, heavy nuclei having magnetic rigidities above the cutoff for Swarthmore, Pennsylvania, presumably reach atmospheric depths attainable by balloons.

If the burst rate at high altitudes were now lower than in 1950, this would be indicative of any or all of the following:

A. The heavy-nucleus flux is reduced at the low-energy end of the spectrum by the cutoff mechanism which is operative on the protons.[1]

B. The intensity of heavy nuclei is reduced throughout the energy spectrum.

C. The intensity of high-energy primary protons is reduced.

Although thus far data are available from only two pulsed-ionization chamber flights conducted during 1958, the results are sufficiently definitive to warrant certain conclusions. The justification for confidence in the validity of the data is illustrated in Fig. 1. It is seen here that the 1958 points are in excellent agreement with the intensity *vs* altitude curve obtained in 1950 (see Table I). It is unlikely that this agreement could be fortuitous. This result indicates that none of the possibilities A, B, or C has occurred.

The contribution to bursts >1 Po-α arising from nuclear disintegrations is significant down to a greater atmospheric depth than that from the heavy-nucleus component. Hence, the statistical uncertainty (±4%) in the comparison of the high-energy primary proton fluxes in 1950

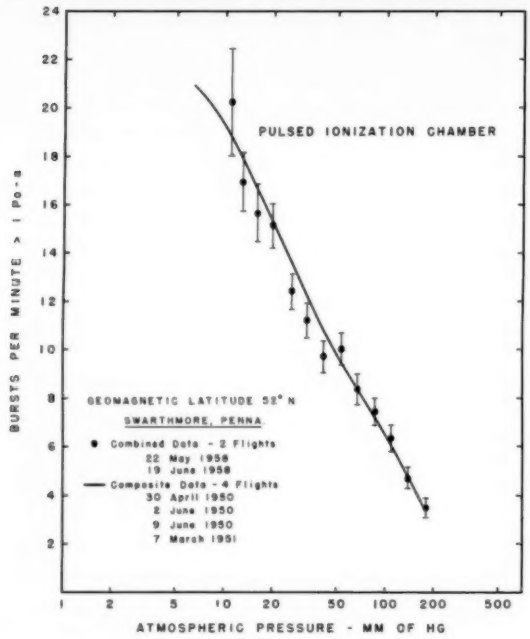

FIG. 1. Pulsed ionization chamber counting rate (bursts >1 Po-α) as a function of atmospheric pressure at Swarthmore, Pennsylvania, in 1950-51 (solid line) and in 1958 (solid dots).

and 1958 is much less than the uncertainty (±13%) in the comparison of the heavy nucleus intensities. Subject to these statistical limitations, the present observations reveal that the fluxes of both heavy nuclei and high-energy protons did not change at a time when the total intensity, as measured by a balloon-borne vertical quadruple-coincidence counter train containing 7.5 cm of Pb, had decreased by ~23%.

Determinations of the α-particle flux and energy spectrum on May 17, 1957 by Freier, Ney, and Fowler[4] had shown that the large reduction in the total intensity of cosmic rays at high altitudes is accompanied by a 45% decrease in the α-particle component, the change occurring throughout the spectrum. Thus, the heavy nuclei in mid-1958 did not behave during a period of reduced total primary cosmic-ray intensity in the same manner as the α-particles a year earlier.

From neutron monitor measurements in aircraft, Meyer and Simpson[5] concluded that the value of γ in the differential spectrum varied from 2.7 to 2.5 between 1954 and 1956, but had not changed between 1951 and 1954. The present results indicate that the current general reduction in the primary cosmic-ray intensity is a consequence of a change in the cutoff rigidity only.

It should be pointed out that a sharp cutoff of protons unaccompanied by a corresponding cutoff of heavy nuclei could be reconciled by ascribing the observed decrease in primary intensity to the action of a cutoff mechanism at the sun upon low-energy cosmic rays of solar origin.

* Assisted by the joint program of the Office of Naval Research and the U. S. Atomic Energy Commission, and by the U. S. National Committee for the International Geophysical Year through the National Science Foundation. Field operations sponsored by the National Geographic Society.

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Table I. Summary of pulsed ionization chamber counting rates for bursts >1 Po-α at Swarthmore, Pennsylvania (geomagnetic latitude 52°N), in 1958 and in 1950-51, respectively. Uncertainties indicated are statistical standard deviations.

Pressure interval (mm of Hg)	Counting rates (bursts/min) May 22, 1958	June 19, 1958	1958 Average	1950-51 average
200-87	4.8±0.3	5.3±0.4	5.0±0.2	5.0±0.2
87-54	7.8±0.5	9.0±0.7	8.2±0.4	8.0±0.3
54-33	10.4±0.6	10.4±0.7	10.4±0.4	10.8±0.4
33-21	12.5±0.6	12.4±0.7	12.5±0.5	13.4±0.6
21-13	16.0±1.6	16.4±0.9	16.3±0.3	18.4±1.3
13-10.7	.....	18.4±1.3	18.4±1.3	18.5±0.8

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PHOTOPRODUCTION OF K MESONS*

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(Received July 3, 1958)

We have studied the reaction $\gamma + p \rightarrow K^+ + \Lambda^0$ using the bremsstrahlung beam of the Cornell synchrotron. The reaction is observed by detecting the K meson and measuring its laboratory angle and momentum. Since the reaction is 2-bodied, these measurements, together with the conservation laws, are sufficient to determine the energy of the incident photon responsible for the production of the K meson. The experimental arrangement is shown in Fig. 1. The photon beam passes through a cylindrical liquid

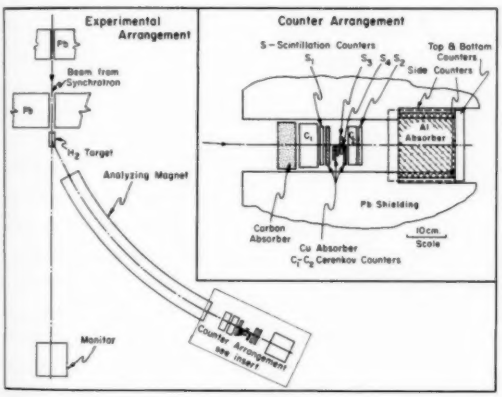

FIG. 1. Experimental arrangement for the detection of K^+ mesons.

hydrogen target (3 in. long and $1\frac{1}{2}$ in. diameter) and is monitored by a total-absorption ion chamber.[1] Positively charged particles emerging from the hydrogen target enter a two-lens strong-focusing analyzing magnet which selects a band of momenta. On emerging from the magnet, the particles pass through a series of counters designed to select the K mesons from the other particles of the same momenta. K mesons pass through the Cerenkov counters C_1 and C_2, the scintillation counters S_1, S_2, and either S_3 or S_4, and come to rest in the aluminum block as is shown in the insert in Fig. 1. They then decay with a mean life of 1.24×10^{-8} sec. About 70% of these give rise to a secondary particle which is detected in one of the four scintillation counters which surround the aluminum block. These four counters are hereafter called "side" counters.

The counters S_3 and S_4 constitute the defining counters for two independent channels (designated 3 and 4) differing in momentum by approximately 13%. A particle is identified as a K meson if it makes a prompt coincidence between counters S_1, S_2, S_3, or S_4, and at least one of the "side" counters. Suitable absorbers are provided so that the pulse heights required in counters S_1, S_2, S_3, and S_4 are about twice that of a relativistic particle. Light particles (e, μ, π) are rejected both by the pulse-height requirement and by C_1 and C_2 in anticoincidence. The K-meson velocity is below the detection threshold in the Cerenkov counters. Particles of proton mass or heavier, of the proper momentum, will not have sufficient range to reach counter S_2. Without the side-counter requirement the background rate is approximately 5 to 10 times the K-meson rate. This background rate, which corresponds to about 0.5% of the π mesons and protons traversing the magnet, probably arises from the particles of the wrong momentum which get through the analyzing magnet in some unorthodox manner, i.e., scattering from the pole tips, and so do not have the ranges and pulse heights predicted for particles of the proper momentum. The addition of the side-counter requirement reduces the background in most cases to less than 10%.

The prompt coincidences, S_1, S_2, S_3, or S_4, plus one of the side counters is made with a resolving time $\tau = 5\times10^{-7}$ sec. In addition, a coincidence is required between S_2 and the side counters with a resolving time $\tau = 5\times10^{-8}$ sec. With this resolving time the combined background from chance coincidences and electrons arising from π-μ-e decay is approximately 2%.

The background of the experiment is measured primarily by operating with the peak beam energy 25 to 50 Mev below that necessary to produce the K mesons under observation. At the laboratory angle of 24.5°, a curve was taken of counting rate as a function of peak beam energy. The momentum for channel 4 was 451 Mev/c and for

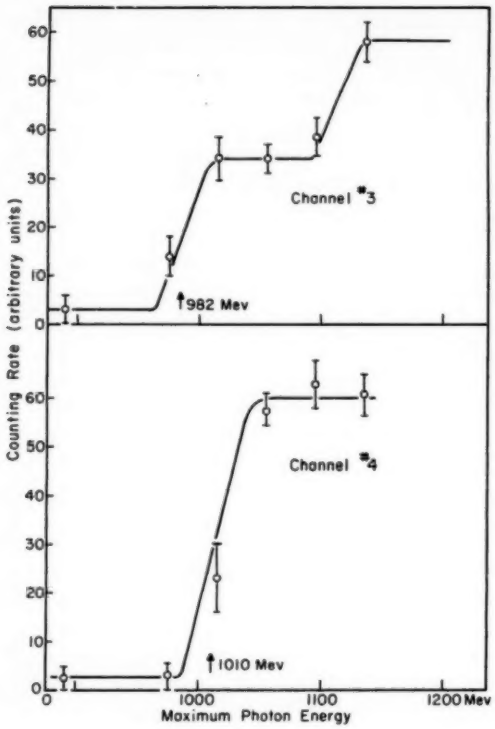

FIG. 2. K^+-meson counting rate as a function of the maximum photon energy of the bremsstrahlung spectrum at the laboratory angle of 24.5 degrees. The two momentum channels were held fixed and sensitive to K mesons produced by photons of 982 and 1010 Mev for channels 3 and 4, respectively.

channel 3 was 398 Mev/c corresponding to K mesons produced by 1010-Mev and 982-Mev photons, respectively. The results are shown for both channels in Fig. 2. The curves drawn are the calculated relative counting rates as a function of peak beam energy, folding in the bremsstrahlung spectrum with the momentum and angular resolution of the magnet. The curves are normalized to the plateau counting rates. In channel 4, one sees a background rate of approximately 5%. The observed relative rates are in good agreement with those calculated. In channel 3, one sees approximately the same behavior except that the point at 1136 Mev lies much higher than the plateau region between 1000 and 1100 Mev. We attribute this increase to K mesons produced in association with Σ^0 hyperons in the reaction $\gamma + p \rightarrow K^+ + \Sigma^0$. The expected relative counting rate for this reaction was calculated in a manner similar to that for the (K^+, Λ^0) reaction. The point at 1136 Mev lies on the plateau of the expected counting rate. The curve in the figure is the calculated rate including the increased K-meson yield resulting from the (K^+, Σ^0) reaction. The background at the lower energies is taken to be a measure of the non-K-particle background. There is also a 3% background of K particles contributed by the walls of the hydrogen target.

The results of the measurement to date are shown in Figs. 3 and 4. Figure 3 shows the cross section, $d\sigma(\theta)/d\Omega$, as a function of the center-of-mass angle, $\bar{\theta}_{c.m.}$, for photons of

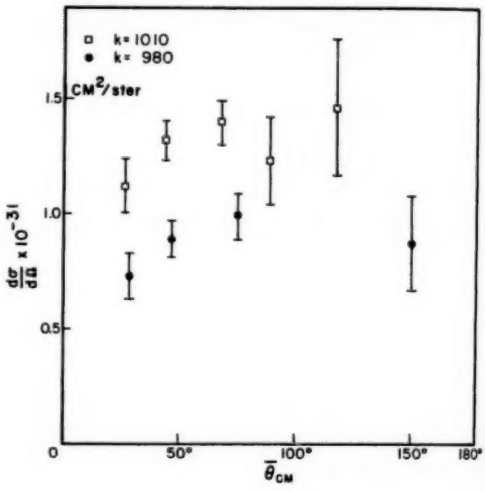

FIG. 3. The cross section for K^+-meson production from hydrogen as a function of center-of-mass angle for photon energies of 1010 and 980 Mev.

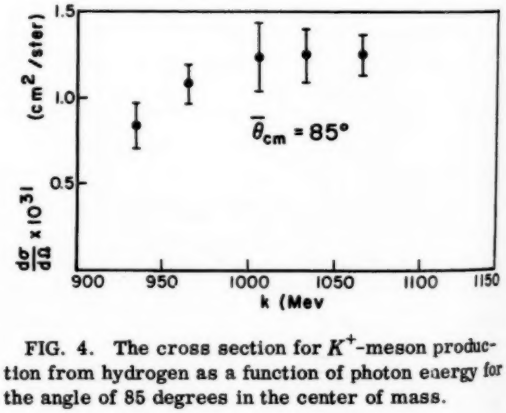

FIG. 4. The cross section for K^+-meson production from hydrogen as a function of photon energy for the angle of 85 degrees in the center of mass.

Table I. A tabulation of the cross sections obtained for the reaction $\gamma+p\rightarrow K^{+}+\Lambda^{0}$. In addition, a single observation is given for the reaction $\gamma+p\rightarrow K^{+}+\Sigma^{0}$. $d\sigma(\theta)/d\Omega$ is given in units of 10^{-31} cm^2/sterad.

$\gamma + p \rightarrow K^+ + \Lambda^0$					
k = 1010 Mev		k = 980 Mev		$\bar{\theta}_{c.m.} = 85°$	
$\bar{\theta}_{c.m.}$	$d\sigma(\theta)/d\Omega$	$\bar{\theta}_{c.m.}$	$d\sigma(\theta)/d\Omega$	k	$d\sigma(\theta)/d\Omega$
26°	1.12 ± 0.12	28°	0.73 ± 0.10	935	0.83 ± 0.16
44°	1.32 ± 0.07	46°	0.90 ± 0.09	963	1.08 ± 0.19
68°	1.40 ± 0.10	75°	0.99 ± 0.10	1006	1.23 ± 0.20
89°	1.26 ± 0.19	150°	0.87 ± 0.21	1032	1.25 ± 0.18
116°	1.46 ± 0.29			1066	1.25 ± 0.14

$\gamma + p \rightarrow K^+ + \Sigma^0$		
k	$\bar{\theta}_{c.m.}$	$d\sigma(\theta)/d\Omega$
1111 Mev	85°	0.61 ± 0.17

energy, k, of 1010 and 980 Mev. The point at 150° in the 980-Mev angular distribution is taken from a measurement previously reported.[2] The scarcity of data and the large probable errors for the measurement backward of 90° center-of-mass angle reflect the low counting rates due to unfavorable kinematical factors, as well as the relatively higher background. The errors shown are the standard statistical errors based entirely on the number of counts. In addition, one should add other uncertainties arising from corrections and possible systematic errors amounting to perhaps 10%. Figure 4 shows the cross section as a function of photon energy for a center-of-mass angle of 85°. Table I shows the numerical results for all measurements including the single measurement made for the associated production of K^+ and Σ^0. Results obtained at California Institute of Technology for the (K^+, Ω^0) reaction are in reasonable agreement with these.[3]

No attempt is made here to compare these results with theoretical calculations that have been published.[4-6] The calculated results for both the scalar and pseudoscalar K mesons are very sensitive to the choice of the anomalous magnetic moments of the proton and the hyperon and it is at present unclear to us what choice is appropriate for these moments. There are a variety of choices of anomalous moments for both the scalar and pseudoscalar calculations which fit the data reasonably well.

We would like to express our gratitude to Mr. Duane Olson for his invaluable aid in the latter stages of this experiment.

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MULTIPLE MESON PRODUCTION IN ELECTRON-PROTON COLLISIONS

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Recently there has been some interest in the possibility of obtaining information on the structure of the proton by observing meson production arising in electron-proton collisions.[1,2] In this note some calculations pertaining to multiple meson production obtained in this process are discussed. We consider the reaction

$$e + p \rightarrow e + p + \pi^{(+)} + \pi^{(-)}, \qquad (1)$$

since the measurement of negative pion production rates are most accessible to experimental

determination.

It is possible to define a "normal distribution" giving the ratios of meson yields in the electron-proton collision process as compared with those obtained in the photoproduction process. This essentially involves the Weizsäcker-Williams approximation of replacing the current matrix element by its transverse part corresponding to an electromagnetic wave propagating in the direction of the incident beam. The method is essentially that outlined in reference 1.

We start with the differential cross section for multiple meson photoproduction:

$$\int_{k_{\min}}^{k_{\max}} (d^2\sigma_k/d\omega_{\vec{q}}\, d\Omega_{\vec{q}})\,(dk/k)$$

$$= (1/128\pi^5) \int\cdots\int (M^2/E_{\vec{p}} E_{\vec{p}'})\, qq'\omega_{\vec{q}}\,\omega_{\vec{q}'}$$

$$\times (dk/k^2)\, d\Omega_{\vec{q}'}\, \Phi_T(\vec{k}, \vec{q}, \vec{q}'). \qquad (2)$$

For the electron-proton collision process, the cross section is

$$(d^2\sigma_p/d\omega_{\vec{q}} d\Omega_{\vec{q}})$$

$$= (1/128\pi^7) \int\cdots\int (M^2/E_{\vec{p}} E_{\vec{p}'})\, dp'\, d\Omega_{p'}\, d\Omega_{q'}$$

$$\times \{qq'\,\omega_{\vec{q}}\,\omega_{\vec{q}'}/(k^2 - k_0^2)^2\}\, \Phi(k, k_0, \vec{q}, \vec{q}'). \qquad (3)$$

The notation here is that of reference 1. We may now substitute the transverse matrix element for current, as given in (2), into relation (3). The result is

$$|dY_p/dY_k| = N_e(p, p')\, k_f, \qquad (4)$$

where

$$Y_p = (d^3\sigma_p/d\omega_{\vec{q}}\, d\Omega_{\vec{q}}\, dp'),\ Y_k = d^2\sigma_k/d\omega_{\vec{q}}\, d\Omega_{\vec{q}},$$

with

$$k_f = p - p_f'.$$

The latter quantities (Y's) can be determined experimentally by observing pion yields as functions of the energies of the incident beams. The right-hand side of (4) is the "normal distribution" defined in reference 1.

Equation (4) may be used as a basis of comparison, since the actual ratio will involve contributions from the longitudinal components of current as well as the transverse parts. The total effect has been computed using the Chew-Low matrix elements for the current.[3,4] As pointed out in reference 1, the static approximation to the longitudinal matrix elements ought

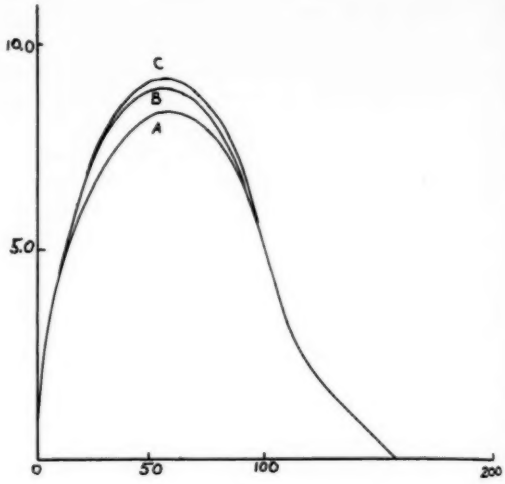

FIG. 1. Theoretical curve of $d^2\sigma_p/d\omega_{\vec{q}} d\Omega_{\vec{q}}$ in micro-microbarns per steradian per Mev as a function of the kinetic energy of the negative pions. Curve (A) represents the result obtained using the transverse matrix element alone. Curve (B) represents the result using both longitudiaal and transverse matrix elements. Curve (C) represents the result obtained if the longitudinal matrix element is used without the factor (5).

to be modified by a factor

$$2k_0 E/(E^2 - M^2). \qquad (5)$$

The results of the calculation are illustrated in Fig. 1 where it may be seen that the effect of the longitudinal current is fairly small. Thus it is not clear whether it is possible to construct an experiment sufficiently sensitive to detect the effect of a factor such as (5).

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BETA-NEUTRINO-ALPHA DIRECTIONAL CORRELATION*

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(Received July 3, 1958)

Recently we have proposed thirty-three possible experiments to decide the relative strength of scalar and vector interactions and that of ten-

sor and axial vector interactions in beta decay.[1] Those are measurements of the recoil nucleus and of the nuclear fluorescence resonance in the decays involving beta-gamma, beta-gamma-gamma, and K capture-gamma-gamma cascades. Now a similar consideration is done for the decay involving a beta-alpha cascade.

The beta-neutrino-alpha directional correlation is immediately derived from the beta-neutrino-gamma angular correlation,[2] replacing $(L_1L_1'\,1\text{-}1\,|n0)$ in $F_n(L_1L_1'j_2j_1)$ by $-(L_1L_1'00\,|n0)$, where both L_1+L_1' and n are even.[3] The result for allowed beta transitions in the decay scheme $j(\beta)j_1(\alpha)j_2$ is

$$W(\vec{p},\vec{q},\vec{k},E)dEd\Omega_e d\Omega_\nu d\Omega_\alpha$$
$$= F(Z,E)pq^2EdEd\Omega_e d\Omega_\nu d\Omega_\alpha\,\xi\left(\left[1+a(\vec{p}\cdot\vec{q}/Eq)+(b/E)\right]\left[\textstyle\sum_{L_1}(j_2\|L_1\|j_1)^2\right]+c'\left[3(\vec{p}\cdot\vec{k})(\vec{q}\cdot\vec{k})-\vec{p}\cdot\vec{q}\right]\right.$$
$$\times(Eq)^{-1}\left\{\textstyle\sum_{L_1,L_1'}(-)^{j_2-j_1}\left[(2j_1+1)(2L_1+1)(2L_1'+1)\right]^{\frac{1}{2}}(L_1L_1'00\,|20)W(j_1j_1L_1L_1';2j_2)\right.$$
$$\left.\left.\times(j_2\|L_1\|j_1)(j_2\|L_1'\|j_1)\right\}\right) \qquad (1)$$

Here we use the same notation as in reference 2, except that $\vec{k}$ is the unit momentum vector of the alpha particle and $(j_2\|L_1\|j_1)$ is the relative intensity of the L_1th partial wave of the alpha particle. c' is given by

$$c'\xi=\left[\left(|C_T|^2+|C_T'|^2-|C_A|^2-|C_A'|^2\right)\pm(\alpha Z/p)\right.$$
$$\times 2\,\mathrm{Im}(C_TC_A^*+C_T'C_A'^*)\left.\right]M_{GT}^2$$
$$\times\left[(2j_1+3)(j_1+1)/45j_1(2j_1-1)\right]^{\frac{1}{2}}\Lambda_{jj_1} \qquad (2)$$

for $\beta^{\mp}$, with

$$\Lambda_{jj_1}=1 \qquad \text{for } j-j_1=j+1,$$
$$=-(2j_1-1)/(j_1+1) \qquad \text{for } j-j_1=j,$$
$$=j_1(2j_1-1)/(2j_1+3)(j_1+1) \qquad \text{for } j-j_1=j-1.$$

The directional correlation function of the beta and alpha particles and the recoil nucleus will be given from Eq. (1) by substituting $1\to 2\pi/q$ and $(\vec{p}\cdot\vec{q})$ and $(\vec{p}\cdot\vec{k})(\vec{q}\cdot\vec{k})\to -2\pi(\vec{p}\cdot\vec{k})(\vec{p}\cdot\vec{k}+\vec{R}\cdot\vec{k}+K)/q$, with $q=|\vec{p}\cdot\vec{k}+\vec{R}\cdot\vec{k}+K|$. $\vec{R}$ and K are the momentum of the nuclear recoil and the absolute value of the momentum of the alpha particle, respectively.

The simplest example of the beta-alpha cascade is $Li^8(\beta)Be^{8*}(\alpha)\,He^4$. In fact, an experiment for this beta decay has been done by Lauterjung, Schimmer, and Maier-Leibnitz.[4] They have counted the beta-alpha-alpha coincidence rate for the geometry indicated in Fig. 1, for which Eq. (1) reduces to

$$W(\vec{p},\vec{q},\vec{k},E)$$
$$= F(Z,E)pq^2E\left[1+\alpha X(\vec{p}\cdot\vec{k}/E)(\vec{q}\cdot\vec{k}/q)\right], \qquad (3)$$

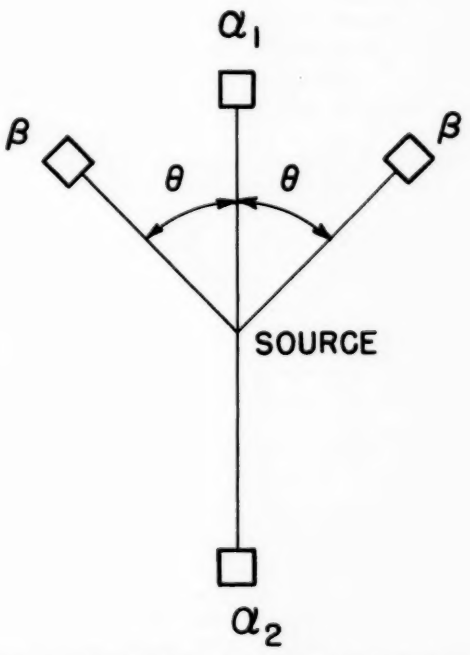

FIG. 1. Geometry for beta-alpha-alpha directional correlation adopted by Lauterjung, Schimmer, and Maier-Leibnitz,[4] and for Eq. (3) in the text. Two alpha counters, α_1, α_2, and two beta ones, β, are placed in a plane. Here α_1 and α_2 are in a line. The two β's have the same angle θ with respect to the direction of α_1. The coincidence rate of α_1, α_2, and β is measured.

with

$$\alpha = -1/3 \text{ for } 1_+ - 2_+ - 0_+,$$
$$= 1 \text{ for } 2_+ - 2_+ - 0_+,$$
$$= 1/7 \text{ for } 3_+ - 2_+ - 0_+,$$

and

$$X = (|C_T|^2 - |C_A|^2)/(|C_T|^2 + |C_A|^2).$$

Here we assume $C_T = -C_T'$, $C_A = C_A'$, and $M_F = 0$. The last assumption comes from the selection rule of the isotopic spin. Therefore, if the decay scheme is $2_+ - 2_+ - 0_+$, this experiment is a more sensitive determination of $|C_T/C_A|^2$ than the electron-neutrino angular correlations of He^6 and Ne^{23}, because the anisotropy for the former is one while that of the latter is only one-third. If M_F is not equal to zero, αX for $2_+ - 2_+ - 0_+$ should be replaced by $[(-|C_S|^2 + |C_V|^2)M_F^2 + (|C_T|^2 - |C_A|^2)M_{GT}^2] / [(|C_S|^2 + |C_V|^2)M_F^2 + (|C_T|^2 + |C_A|^2)M_{GT}^2]$. In the case of $2_+ - 2_+ - 0_+$, Eq. (3) coincides with Eq. (1) in reference 4. Since the value of $d = \alpha X$ is large and negative with certain experimental error,[4] we can conclude not only that the axial vector is predominant, but also that the spin of Li^8 is 2+ and not 1+ or 3+.

Formulas for more complicated decay schemes involving alpha, beta, and gamma transitions will be derived easily from several formulas in reference 1 with the rule in reference 3.

The author would like to express his sincere thanks to Professor K. H. Lauterjung for a valuable discussion at Columbia University.

*Work partially supported by the U. S. Atomic Energy Commission.

[1] Morita, Morita, and Yamada, Phys. Rev. 111 237 (1958). M. Morita and R. S. Morita, Bull Am. Phys. Soc. Ser. II, 3, 270 (1958), and Phys. Rev. (to be published).

[2] M. Morita, Nuclear Phys. 6, 132 (1958).

[3] The beta-alpha directional correlation for the forbidden transitions was given by M. Morita, Progr. Theoret. Phys. Japan 15, 445 (1956), especially p. 456, and M. Morita and R. S. Morita, Phys Rev. 110, 461 (1958), especially p. 464. That for the allowed beta transitions was given by M. Morita and M. Yamada, Progr. Theoret. Phys. Japan 13, 114 (1955), taning into account higher order corrections.

[4] Lauterjung, Schimmer, and Maier-Leibnitz, a post-deadline paper presented at the Ithaca Meeting of the American Physical Society, 1958, and Z. Physik (to be published).

LARGE-ANGLE ELECTRON PAIR PRODUCTION FROM HYDROGEN*

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(Received June 26, 1958)

In this experiment, 70-Mev positrons produced in a liquid hydrogen target are detected at 90° to a 137-Mev peak energy bremsstrahlung beam made by passing the electron beam from the Stanford Mark III linear accelerator through a tantalum radiator. Figure 1 shows the two main Feynman diagrams for the process. It is important to note here that uncertainties in the pair-production calculation due to the influence of proton structure on the rescattering of the virtual intermediate lepton can be uniquely removed from the problem. The results of the electron-proton scattering experiments can be used to predict the effect of the proton in pair production. The e-p scattering form factors depend only on the four-momentum transfer and can be introduced phenomenologically into the pair-production calculation. This leaves the electrodynamics as the only source of error in the calculation. A complete calculation of the pair production cross section from protons with form factors including radiative corrections and Compton terms has been carried out by Bjorken, Drell, and Frautschi.[1]

In Fig. 1(a) a photon of momentum k produces a pair, and a positron of momentum p_+ comes directly out. For the conditions of this experiment the four-momentum transfer q at the pair vertex is given by $q^2 = 2kp_+$. In Fig. 1(b) the electron comes directly out and q is of the order of $m_e c$. The mean value of q for diagram 1(a)

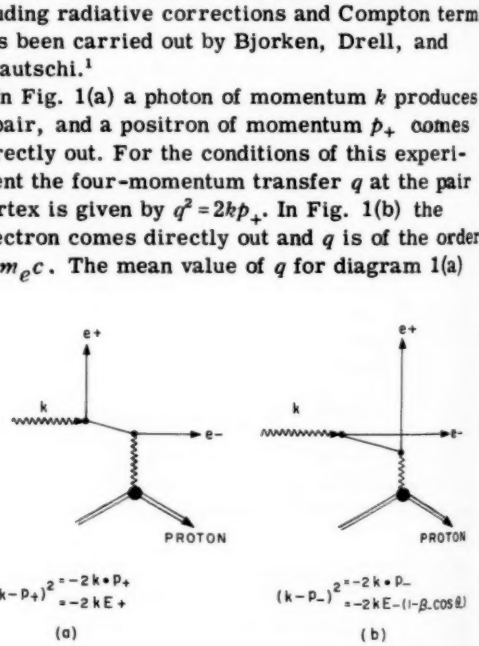

FIG. 1. Principal Feyman diagrams for pair production from a proton.

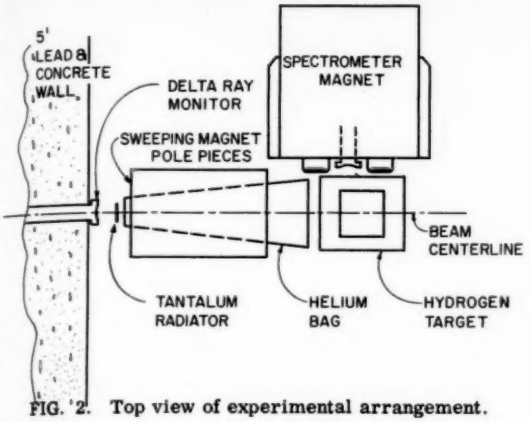

FIG. 2. Top view of experimental arrangement.

is ~ 115 Mev/c, making the experiment useful as a test of the small-distance (large-momentum-transfer) behavior of quantum electrodynamics. Diagram 1(a) contributes about 30% to the cross section.

Figure 2 shows a top view of the experimental setup. The energy-analyzed electron beam passes through a delta-ray monitor, a 0.1-radiation-length tantalum converter, and a magnet which sweeps charged particles from the beam, and then hits a Styrofoam liquid hydrogen target. The positrons are momentum-analyzed by a double-focusing magnetic spectrometer and detected by two Plexiglas Cerenkov counters in coincidence.

The main source of background in the experiment comes from events wherein a pair is produced either in the target walls or in the hydrogen, and the positron is then scattered into the spectrometer by a nucleus other than the one which produced it. About 50% of the counting rate comes from this source. The data were normalized by measuring the yield of elastically scattered electrons from the liquid hydrogen target using exactly the same setup as was used in measuring the positron yield, and using the measurements of Hofstadter and co-workers[2] of the proton form factor in the Rosenbluth cross section to determine the normalization factor.

The result to date is that the ratio of the experiment to the prediction of Bjorken, Drell, and Frautschi is 0.96 ± 0.14 where most of the error comes from counting statistics. The question now is, What does it mean? There is no good way to characterize a deviation from the point interaction theory of quantum electrodynamics without having constructed another theory to replace it. Lacking a theory, all one can do is make a guess. If the electron propagator is modified,[3]

$$(p^2-m_0^2c^2)^{-1} \rightarrow \{(p^2-m_0^2c^2)^{-1} - [p^2-m_0^2c^2-(\hbar^2/\Lambda_e)]^{-1}\},$$

one obtains a form factor for the electrodynamic interaction that looks like $F^2(q^2) = 1 - 2q^2\Lambda_e^2 + \ldots$.

This expression has no fundamental theoretical significance: it is merely a way to characterize a deviation from the point interaction theory. Figure 3 shows a plot of the cutoff distance vs the ratio of the yield predicted by a theory with cutoff to that by the point-interaction theory. It has been assumed that diagram 1(b), which contributes 70% of the cross section but involves only a small momentum transfer, will agree with the point theory. The result of this experiment is indicated on Fig. 3. The experimental error (standard deviation) intersects the curve at a value of Λ_e of about 0.9 fermi. This might be an upper limit to the distance at which the present theory of quantum electrodynamics breaks down.

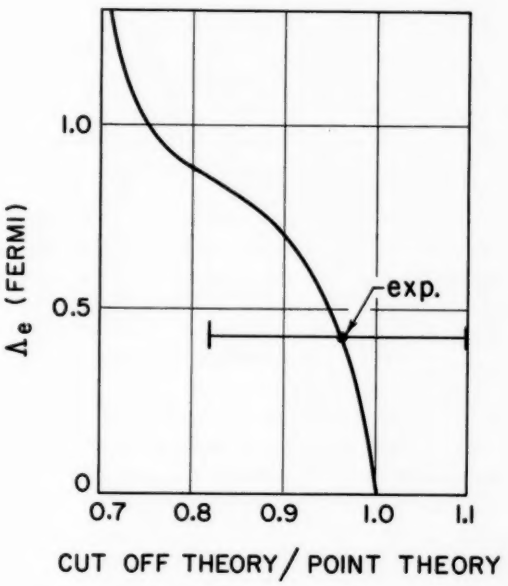

FIG. 3. Cutoff parameter as a function of the ratio of the predicted yield of a cutoff theory to that of the point interaction theory for the conditions of this experiment.

* Supported by the joint program of the Office of Naval Research, U. S. Atomic Energy Commission, and Air Force Office of Scientific Research.

[1] Bjorken, Drell, and Frautschi, Phys. Rev. (to be

published).

[2] Hofstadter, Bumiller, and Yearian, Revs. Modern Phys. (to be published).

[3] For definition of the notation, and for a general discussion of this and other aspects of quantum electrodynamics at small distances, see S. D. Drell, Ann. Phys. 4, 75 (1958).

POLARIZATION OF CONVERSION ELECTRONS FOLLOWING β DECAY

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(Received May 23, 1958;
revised manuscript received June 19, 1958)

The experiments (aside from the use of polarized sources) which serve to verify breakdown of parity and charge conjugation invariance in the β interaction are largely of two types: measurement of the longitudinal polarization of the β particles, and of the circular polarization of γ radiation following the β transition.[1] In addition to the demonstration of P and C breakdown, these experiments are useful as a method of testing the validity of the two-component theory. In addition, variations of this method, involving observation of the recoil could, in principle, provide information relevant to the β interaction.[2] The β-particle polarization analysis entails conversion of longitudinal to transverse polarization and the consequent loss of intensity. The β-circularly polarized γ correlation is similarly handicapped by the circumstance the only two (approximately) of the Fe electrons contribute to the magnetization of the analyzing field.

In this communication we discuss an alternative procedure which may offer some advantages particularly for β transitions followed by moderately low-energy conversion electrons. This involves the measurement of the polarization of the conversion electrons in coincidence with the β. The primary point of interest here is that these electrons are partially transversely polarized and that in many cases the polarization is fairly large. The existence of a transverse polarization in an allowed transition is made possible by the coincidence observation, defining a plane in which the polarization lies, and by the finite rest mass of the conversion electrons. These are to be understood as necessary conditions.

The (unnormalized) population of the nuclear state resulting from β decay referred to the direction of the β particle is

$$p_M = 1 + bM. \tag{1}$$

The present results apply to any process wherein a polarization of the form (1) exists.[3] If $\vec{n}'$ is a unit vector defining the direction of the polarization in the rest system, the polarization in the laboratory frame is

$$P_n = (\Psi, S_n \Psi)/(\Psi, \Psi), \tag{2}$$

where[4]

$$S_n = i\gamma_5\gamma_\mu n_\mu = -\beta\vec{\sigma}\cdot\vec{n} - \beta\gamma_5\vec{n}\cdot\vec{p}/E. \tag{3}$$

Here Ψ is the modified plane wave for the conversion,[5] n_μ is obtained from $\vec{n}'$ by a Lorentz transformation and $\vec{p}$, E are the momentum, total energy of the conversion electron ($m, c = 1$). For transverse polarization $S_n = -\beta\vec{\sigma}\cdot\vec{n}$ and $\vec{n} = \vec{r}_1 \times (\vec{e}\times\vec{r}_1)/\sin\theta$ with $\vec{r}_1$ a unit vector along $\vec{p}$, $\vec{e}$ a unit vector in the direction of the β, and θ the angle between these two directions. For longitudinal polarization $\vec{n}p = E\vec{p}$, and S_n is equivalent to $\vec{\sigma}\cdot\vec{p}/p = \vec{\sigma}\cdot\vec{r}_1$.

The decay chain is defined by $J_1(\beta)\, J(c.e.)\, J_2$. The multipolarity of the conversion link is most generally a magnetic L - electric L' mixture and $L' = L \pm 1$ (rigorously). With these definitions the polarization can be written as ($q = \parallel$ or $\perp$)

$$P_q = \frac{P_q(ML) + \bar{\delta}^2 P_q(EL') + 2\bar{\delta}\, P_q(ML, EL')}{1 + \bar{\delta}^2}, \tag{4}$$

where $\bar{\delta}^2$ is the ratio of intensities of EL' to ML conversion electrons:

$$\bar{\delta} = \delta(\alpha_{L'}/\beta_L)^{\frac{1}{2}}, \tag{5}$$

where δ is the parameter[6] whose square gives the relative γ-ray intensities EL' to ML and $\alpha_{L'}$ and β_L are, respectively, the electric and magnetic conversion coefficients for pure multipoles.

Then for the pure multipole longitudinal polarization one finds, for an initial $s_{\frac{1}{2}}$ electron,

$$P_\parallel(ML) = \tfrac{1}{2} bN\cos\theta \times \frac{|Q_{L+1}\exp(i\eta_{L+1}) - Q_{-L}\exp(i\eta_{-L})|^2}{(L+1)|Q_{L+1}|^2 + L|Q_{-L}|^2}, \tag{6}$$

$$N = \frac{L(L+1) + J(J+1) - J_2(J_2+1)}{2L+1}. \tag{6a}$$

$\eta_\kappa\,(\kappa = L+1, -L)$ are the (Coulomb) phase shifts and the Q_κ are defined in reference 5. For co-

incidences with $\beta^{\pm}$ particles of velocity v,

$$\xi b=\frac{2v}{c}\ \mathrm{Re}\Big\{\mp(C_A C_{A'}{}^{*}-C_T C_{T'}{}^{*})\lambda_{J_1 J}|M_{GT}|^2 + (C_S C_{T'}{}^{*}+C_{S'}C_T{}^{*}-C_V C_{A'}{}^{*}-C_{V'}C_A{}^{*}) \times \delta_{J_1 J}[J_1(J_1+1)]^{-\frac{1}{2}}M_F M_{GT}{}^{*}\Big\}, \quad (6b)$$

$$\xi=\Big\{|C_S|^2+|C_{S'}|^2+|C_V|^2+|C_{V'}|^2\Big\}|M_F|^2 +\Big\{|C_T|^2+|C_{T'}|^2+|C_A|^2+|C_{A'}|^2\Big\}|M_{GT}|^2,$$

$$\lambda_{J_1 J}=(J_1+1)^{-1}, \quad J=J_1+1,$$
$$=[J_1(J_1+1)]^{-1}, \quad J=J_1,$$
$$=-1/J_1, \quad J=J_1-1$$

in the standard notation. For pure electric transitions,

$$P_{\parallel}(EL)=\tfrac{1}{2}bN\cos\theta \times\frac{|Q_{-L-1}\exp(i\eta_{-L-1})-Q_L\exp(i\eta_L)|^2}{(L+1)|Q_{-L-1}|^2+L|Q_L|^2}, \quad (7)$$

and the Q_κ ($\kappa=-L-1, L$) now refer to electric multipoles.

For the pure transverse polarizations the results are

$$P_{\perp}(ML)=-\tfrac{1}{2}bN\sin\theta\ \frac{(L+1)|Q_{L+1}|^2-L|Q_{-L}|^2-\mathrm{Re}\{Q_{L+1}Q_{-L}{}^{*}\exp[i(\eta_{L+1}-\eta_{-L})]\}}{(L+1)|Q_{L+1}|^2+L|Q_{-L}|^2}, \quad (8)$$

$$P_{\perp}(EL)=\tfrac{1}{2}bN\sin\theta\ \frac{(L+1)|Q_{-L-1}|^2-L|Q_L|^2-\mathrm{Re}\{Q_{-L-1}Q_L{}^{*}\exp[i(\eta_{-L-1}-\eta_L)]\}}{(L+1)|Q_{-L-1}|^2+L|Q_L|^2}. \quad (9)$$

The mixture terms have also been calculated but will be omitted here in the interest of brevity. They will be given together with complete numerical results for all pure and mixed terms in a later publication where forbidden transitions will also be discussed.

It will be noted that the electric monopole contribution is strictly zero and that there are no interference terms of the $E0 + E2$ type. In fact, $|L-L'|=1$ for interference between any pair of multipoles. The sign of $P_q(ML)$ and $P_q(EL)$ can be inferred very simply for the longitidinal case, and for the magnetic transverse polarization the sign and approximate magnitude of $P_\perp$ is readily deduced from the fact that in most cases $Q_{-L} \gg Q_{L+1}$. In the limit $E\to\infty$, the $P_\perp$ goes to 0 and $P_\parallel$ goes to circular polarization of the competing γ ray, as expected. In the nonrelativistic limit $P_q(EL)$ goes to 0 also as expected.

As an indication of the magnitude of the polarization we give one numerical example. We consider the β^- emitter I^{131} and in particular the $E2$ (284-kev) transition to the first excited state of the daughter Xe^{131} following the main (87.2%) branch. The spins are $J_1=7/2$, $J=5/2$, $J_2=1/2$, and all parities are even.[7] At $\theta=\frac{1}{2}\pi$ we find for K conversion[8] and with the two-component theory $P_\perp(E2)=-0.34\ v/c$ and at $\theta=\pi$, $P_\parallel(E2)=0.13\ v/c$. The signs of P_q imply a polarization in the direction opposite to the direction of the β^- in both cases. For the 364-kev ($E2+M1$) transition to ground, $P_\perp=0.28\ v/c$ at $\theta=\frac{1}{2}\pi$. Other sources exist for which $P_\perp$ is larger but I^{131} seems to be particularly suitable for many other reasons, e.g., convenient lifetime of the transitions involved.

The authors are indebted to Dr. R. R. Lewis for a very valuable discussion.

[1] These subjects have been reviewed by H. Frauenfelder and C. S. Wu and by R. M. Steffen, Proceedings of the Rehovoth Conference on Nuclear Structure (North-Holland Publishing Company, Amsterdam, 1958).

[2] S. B. Treiman (to be published). Also Frauenfelder, Jackson, and Wyld (to be published). This work was motivated by a remark of the latter authors who anticipated, in a qualitative way, the longitudinal part of the polarization discussed here.

[3] For the recoil problem in $\beta^{\pm}$ emission the nuclear state also exhibits alignment (reference 2) arising from the GT part of the interaction. However, for recoil in K capture, p_M has the form (1).

[4] H. A. Tolhoek, Revs. Modern Phys. 28, 277 (1956); see appendix, p. 297.

[5] Rose, Beidenharn, and Arfken, Phys. Rev. 85, 5 (1952); see Eq. (21) of this paper.

[6] L. C. Biedenharn and M. E. Rose, Revs. Modern Phys. 25, 729 (1953); see Sec. III. A of this paper.

[7] Strominger, Hollander, and Seaborg, University of California Radiation Laboratory Report UCRL-1928,

April, 1958 (unpublished).

[8] All numerical results are obtained from a tabulation of conversion matrix elements which was made for the K shell, point nuclei and no screening. This tabulation will appear in a forthcoming publication, M. E. Rose, Internal Conversion Coefficients, [North-Holland Publishing Company, Amsterdam (to be published)]. Since in the case considered $Q_{-L-1} >> Q_L$ the results are almost entirely unaffected by details of the electron dynamics.

ERRATA

The following Letter originally appeared in Phys. Rev. Lett. 1, 52 (1958) (see Editorial, page 90).

POLARIZATION OF FREE POTASSIUM ATOMS BY EXCHANGE COLLISIONS WITH SODIUM ATOMS AND FREE ELECTRONS*

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(Received June 9, 1958)

We wish to report an experiment in which free potassium atoms have been detectably polarized by exchange collisions with free polarized sodium atoms[1] and with free polarized electrons. The apparatus is very similar to that discussed by Dehmelt[2] for his electron polarization experiment in which free electrons were polarized by exchange collisions with an optically pumped (polarized) sodium vapor.

Our sample consists of a 6-cm diameter Pyrex bulb into which has been distilled a mixture of sodium and potassium metal,[3] together with 6 mm argon. There are two 0.03-in. tungsten wire electrodes oppositely located in the bulb, which are employed in creating momentary discharges in the sample. The bulb is normally held at 170°C at which temperature we estimate, from the optical data, that the density of Na atoms is $\sim 5\times10^{10}$/cc and the density of K atoms is $\sim 2\times10^{10}$/cc.

Na resonance radiation from an Osram spectroscopic lamp is sent through a Na-pass optical filter[4] followed by a circular polarizer. This light passes through the sample and is collected by a photocell, the signal from which is amplified and displayed on a CRO. The direction of light propagation coincides with that of a 9-gauss uniform magnetic field developed by 25 in. diameter Helmholtz coils. A 0.04-gauss peak-to-peak sinusoidal field modulation at 15 cycles/sec is supplied by coils coaxial with the main magnetic field. Radio-frequency magnetic fields are developed at the sample by means of 4 turns of wire wrapped around the bulb and excited by a General Radio Unit Oscillator.

The six Na transitions ($F=2$, $\Delta m=1$), ($F=1$, $\Delta m=1$) are well resolved. An example is shown in Fig. 1(b).

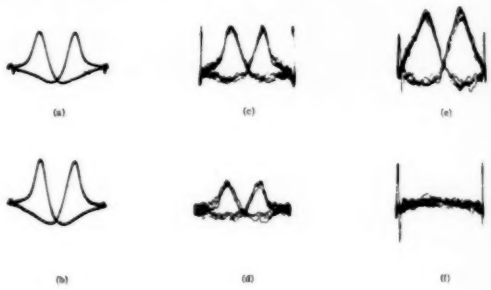

FIG. 1. CRO traces of the sodium, potassium, and free electron resonances. The homogeneous magnetic field is 8.9 gauss, and the modulation field is 0.04 gauss peak-to-peak at 15 cycles/sec. The Na transition ($F=2$; $m=-1 \rightarrow m=-2$) occurs at 6.2 Mc, the K transition ($F=2$; $m=-1 \rightarrow m=-2$) occurs at 6.4 Mc, and the free electron resonance occurs at 25 Mc. The discharge is at 30 cycles/sec and is seen as pips at the extreme ends of the traces.

(a) Na resonance, discharge on. CRO gain = 1 (arbitrary units).
(b) Na resonance, discharge off. Gain = 1.
(c) K resonance, discharge on. Gain = 10.
(d) K resonance, discharge off. Gain = 10.
(e) Free electron resonance, discharge on. Gain = 10.
(f) No radio-frequency, discharge on. Gain = 10.

The corresponding six potassium resonances occur at their proper frequencies with intensities approximately 1/10 that of the Na resonances. An example is shown in Fig. 1 (d). The interpretation is that free K atoms have been polarized by exchange collisions with the optically pumped Na atoms. The rf depolarization of the K atoms is "passed on" to the Na atoms via the exchange process, so that a decrease in the Na light is observed. From a measured alkali relaxation time of $\sim 10^{-2}$ sec and the data given above we compute that the cross section for Na-K exchange collisions is $\sim 5\times10^{-14}$cm^2. This estimate is considered reliable to within a factor of three.

The calculation of Purcell and Field[5] for H-H exchange collisions was extended to the Na-K process, utilizing the parameters of the Na-K antisymmetric attraction given by Herzberg.[6]

The result is $\sim 7\times10^{-14}$cm^2 and is considered reliable to within a factor of two.

Free electrons are introduced into the bulb by initiating 2-Mc/sec rf discharges of ~0.002 sec duration, repetitive at 30 cycles per second, and phased so that the discharges occur at the extreme values of the 15-cycle/sec modulation field. (In the CRO displays of Fig. 1 the discharges occur at the extreme right and left hand sides of the trace.) The Na, K, and free electron resonances are shown in Figs. 1 (a), 1 (c), 1 (e). The interpretation of the free electron resonance, following Dehmelt, is that the free electrons remaining after the discharge become polarized by exchange collisions with the optically pumped Na atoms.

When the discharge is turned on, the sodium resonance decreases by ~20%, while the potassium resonance increases by ~40% (see Fig. 1). We believe the decrease in sodium signal is due to two effects: (1) the discharge tends to "etch out" the sodium vapor, and (2) the free electrons exert a depolarizing effect on the sodium atoms. The potassium resonance increases, however, because the free electrons, polarized by the Na atoms, contribute constructively to the polarization of the potassium atoms via the electron-potassium exchange collisions. The potassium vapor is also etched out by the discharge, but the advantageous electron-potassium exchange effect is generally predominant. The alkali densities and hence the relative changes in the alkali resonances are temperature dependent.

The roles played by the Na and K atoms were reversed by running at higher temperatures (~200°C) and using a potassium arc. The observations are similar to those described above and are in agreement with the computed Na-K exchange cross section.

From the data we estimate that the electron-potassium exchange cross section is $<3\times10^{-14}$ cm^2, in agreement with Dehmelt's estimate for the electron-sodium cross section.[2] We have also obtained this value from unpublished data previously taken with an apparatus almost identical to Dehmelt's, but with potassium vapor only and a potassium arc.

These techniques may be useful in polarizing atomic and ionic species hitherto intractable to direct optical pumping. Experiments along this line are in progress.

*This work has been supported by the U. S. Atomic Energy Commission.

[1] While preparing this manuscript for publication we learned that R. Novick, University of Illinois, has been developing an atom-atom exchange polarization experiment in a sodium-rubidium system; see R. Novick and H. E. Peters, following Letter [Phys. Rev. Lett. 1, 56 (1958)]. We wish to express our gratitude to Professor Novick for several interesting telephone conversations in which ideas common to both of these experiments were discussed.

[2] H. G. Dehmelt, Phys. Rev. 109, 381 (1958).

[3] The Na metal used (Cenco) contains a trace of potassium impurity so that a pre-mixing of the two metals was not necessary.

[4] This filter consists of Corning CS3-69 and CS4-97 plates. It is used to absorb the small amount of potassium resonance radiation emitted by the Osram sodium lamp.

[5] E. M. Purcell and G. B. Field, Astrophys. J. 124, 542 (1956).

[6] G. Herzberg, Spectra of Diatomic Molecules (D. Van Nostrand Company, New York, 1955).

INSTABILITY, TURBULENCE, AND CONDUCTIVITY IN CURRENT-CARRYING PLASMA. O. Buneman [Phys. Rev. Lett. 1, 8 (1958)].

In the fourth line of paragraph 2, the equation $\omega_{pi} = (e\rho/M\epsilon/_0)^{\frac{1}{2}}$ should read $\omega_{pi} = (e\rho/M\epsilon_0)^{\frac{1}{2}}$. In the first display equation, the second term on the left-hand side should read $\omega_{pe}{}^2/(\beta u-\omega)^2$. In the sixth line after the table, the expression $18\,\omega_{pe}$ should read $18/\omega_{pe}$.

DIVERGENCELESS CURRENTS AND K-MESON DECAY. S. Weinberg, R. E. Marshak, S. Okubo, E. C. G. Sudarshan, and W. B. Teutsch [Phys. Rev. Lett. 1, 25 (1958)].

In Eq. (1), "α" should be replaced everywhere by "γ". In Eq. (3), "α" should be replaced everywhere by "ϑ".

ANGULAR DISTRIBUTIONS OF TRITONS FROM THE $C^{14}(d,t)C^{13}$ REACTION. W. E. Moore, J. N. McGruer and A. I. Hamburger [Phys. Rev. Lett. 1, 29 (1958)].

Line 21 reads: "M is the nuclear mass----"; but should read: "$\underline{M}$ is the reduced nucleon mass----".

CONFIGURATION MIXING IN THE C^{14} GROUND STATE. E. Baranger and S. Meshkov [Phys. Rev. Lett. 1, 30 (1958)].

Equation (4) reads "$\delta^2/\alpha^2 = 0.14$"; but should

read "$\gamma^2/\alpha^2 = 0.14$."

Reference 6 reads "Harwell Report T/R-1278"; but should read "Harwell Report T/R-1289."

Reference 9 reads "Proc. Phys. Soc. (London) A210, 534 (1951)"; but should read "Proc. Roy. Soc. (London) A210, 534 (1951)."

ABSTRACTS

In this section are printed the abstracts of Articles that have been forwarded to The American Institute of Physics for publication in THE PHYSICAL REVIEW. In quoting information obtained from this section before the appearance of the corresponding Article, reference should be made to "Physical Review (to be published)" rather than to this Journal.

CATHODE SHEATH IN A LOW-DENSITY GAS DISCHARGE. P. L. Auer, H. Hurwitz, Jr., and S. Tamor, General Electric Research Laboratory, The Knolls, Schenectady, New York (Received April 23, 1958).

The space-charge equations of a low-density gas discharge are derived and applied in the present report to the study of potential distributions in the cathode sheath. In order to facilitate mathematical computation, the contribution of secondary electrons and charge exchange to the potential distribution is neglected. The space charge is generated, consequently, by the primary electrons of the hot cathode and the ions resulting from volume ionization of the gas in the sheath. The effect of arbitrary positive ion currents issuing from the positive column to the sheath is also examined.

A self-consistent method of solution of the space-charge equations is employed in the present investigation. The method essentially involves finding a cross section which corresponds to a particular potential distribution. A number of cross sections of physical significance can be obtained by this method. The results of our numerical calculations allow us to draw certain semi-quantitative conclusions about the mode of evolution of a cathode sheath and positive column with increasing gas pressure and applied voltage. An important conclusion of our calculations is that a positive column will begin to evolve when the number of ionization mean free paths in the discharge gap is of the same order as the square root of the ratio of electron to positive ion mass.

EXISTENCE OF POSITRONIUM IN SOLIDS. Mirza A. Baqi Bég, University of Pittsburgh, Pittsburgh, Pennsylvania (Received March 6, 1958; revised manuscript received May 21, 1958).

It is shown, in connection with a recent experiment, that it is impossible, in principle, to detect excited states of positronium within a solid or diffusing out of a solid. An experiment is proposed which would decide directly whether ground state positronium may be formed in a solid.

MOBILITY OF ELECTRONS AND HOLES IN PbS, PbSe, AND PbTe BETWEEN ROOM TEMPERATURE AND 4.2°K. Robert Allgaier, U. S. Naval Ordnance Laboratory, White Oak, Maryland, and University of Maryland, College Park, Maryland, and Wayne W. Scanlon, U. S. Naval Ordnance Laboratory, White Oak, Maryland (Received April 24, 1958).

Hall coefficient and resistivity measurements have been made on 29 single crystals (mostly synthetic) of PbS, PbSe, and PbTe between room temperature and 4.2°K. Almost all of the samples had extrinsic carrier concentrations of the order of 10^{18} per cm^3, as deduced from the Hall coefficients which were essentially constant over the entire temperature range investigated. Hall mobilities were calculated from the Hall and resistivity data, and were found to increase rapidly with decreasing temperature. Between room temperature and about 50°K the mobility behavior was essentially intrinsic and varied approximately as $T^{-2.2}$. Below 50°K the mobility curves turned gradually toward the horizontal in a manner resembling the residual resistance phenomenon observed in metals. Values as high as 800 000 cm^2/v-sec were attained at 4.2°K despite the large carrier concentrations present in all the samples. A simple experiment was performed which suggests that dislocations are the principal scattering mechanism below 50°K, rather than the charged point defects associated with the extrinsic carriers. The possibility that a high static dielectric constant could explain the large mobilities at low temperatures is discussed.

ENERGY LOSS OF FAST CHARGED PARTICLES IN LiH. Werner Brandt, Radiation Physics Laboratory, Engineering Dept., E. I. duPont de Nemours and Company, Incorporated, Wilmington, Delaware (Received April 2, 1958).

The quantum theory of the electromagnetic interaction of particles and fields in dense matter predicts that fast charged particles lose their kinetic energy to longitudinal electronic excitations influenced by one-electron and collective effects. Both effects are to be expected in metals and insulators, but to varying degrees. A clear identification of collective excitations in insulators is still lacking.

LiH is the best suited crystal for predicting

the one-electron and collective effects on the longitudinal electronic excitations in an insulator, and for verifying such prediction experimentally. The characteristic energy loss spectrum and the mean excitation potential of the LiH crystal are estimated. One finds that the longitudinal excitations in this insulator crystal are of predominantly collective character. The results are so significantly different from those obtained by neglecting collective effects that stopping experiments in single as well as in multiple scattering can clearly identify the collective effects and confirm the theoretical predictions.

SEEBECK EFFECT FLUCTUATIONS IN GERMANIUM. James J. Brophy, Armour Research Foundation of Illinois Institute of Technology, Technology Center, Chicago, Illinois (Received January 9, 1958).

Fluctuations in the Seebeck effect or thermoelectric power of germanium single crystals have been observed. The Seebeck noise-power spectrum varies as reciprocal frequency and may be quantatively predicted from current noise measurements. Carrier density fluctuations responsible for $1/f$ noise are also the source of Seebeck noise. The present results indicate that carrier fluctuations having a $1/f$ spectrum persist even in the absence of net dc current flow. Seebeck noise is also observed in low-noise specimens having nonohmic, noisy electrodes.

OPTICAL ABSORPTION OF Cl_2^- MOLECULE-IONS IN IRRADIATED POTASSIUM CHLORIDE. C. J. Delbecq, B. Smaller, and P. H. Yuster, Argonne National Laboratory, Lemont, Illinois (Received April 16, 1958).

Pure KCl, KCl-Ag, KCl-Tl, and KCl-Pb, after x-ray irradiation at liquid nitrogen temperature, show identical Cl_2^- paramagnetic resonance spectra. The optical absorption of the Cl_2^- molecule-ion has been identified and is found to consist of two bands, one at 365 mμ and one at 750 mμ. This identification was made possible by bleaching irradiated samples with polarized light and observing the effects on the optical and also the paramagnetic resonance spectra. It was possible to reorient the Cl_2^- molecule ions such that a large fraction of them were in only one of the six possible $\langle 110 \rangle$ directions.

CRYSTAL STRUCTURE AND FERRI-MAGNETISM IN $NiMnO_3$ AND $CoMnO_3$. W. H. Cloud, Central Research Department, Experimental Station, E. I. du Pont de Nemours and Company, Wilmington, Delaware (Received April 21, 1958).

X-ray diffraction studies of individual crystals of $NiMnO_3$ and $CoMnO_3$ show that these compounds have the structure of ilmenite (space group $R\bar{3}$). The crystals have a plate-like crystal habit with the three-fold symmetry axis perpendicular to the faces of the platelet. Positions of atoms within the unit cell have been determined by comparing calculated structure factors with those obtained from x-ray powder patterns. There are two Ni-O-Mn configurations that appear to be favorable to a magnetic superexchange interaction, and these are sufficient to extend magnetic ordering in three dimensions.

KINETIC EJECTION OF ELECTRONS FROM W BY Cs AND Li IONS. Paul M. Waters,* The Department of Physics, The University of California, Berkeley, California (Received April 16, 1958).

Total yield and kinetic-energy distribution have been measured for electrons ejected from both atomically clean and gas-covered tungsten surfaces by impact of Li^+ and Cs^+ in the kinetic energy range 150 to 1500 ev. The results indicate that the probability of kinetic emission becomes vanishingly small, $\sim 10^{-6}$, in the vicinity of 150 ev ion kinetic energy and increases to 6.2×10^{-2} and 1.2×10^{-2} at 1500 ev for Li^+ and Cs^+ respectively when the tungsten surface is atomically clean. The adsorption of a monolayer of either O_2 or N_2 increases the efficiency of this process by at least an order of magnitude below 350 ev ion energy, and by a factor of 2 or 3 at energies above some 500 ev. In addition, it was observed that the energy distribution of the electrons ejected does not change significantly as the result of such adsorption. An attempt to measure the probability of formation of Li^- from Li^+ striking an atomically clean tungsten surface showed that such a process, if it occurs at all, occurs with a probability less than about 2×10^{-8} for Li^+ incident with 150 to 1500 ev kinetic energy, while negative ions of undetermined constituency were observed to be formed with a probability of about 2×10^{-3} under the same conditions if the tungsten surface was thoroughly contaminated.

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MAGNETIC SUSCEPTIBILITY OF COPPER-NICKEL AND SILVER-PALLADIUM ALLOYS AT LOW TEMPERATURES. E. W. Pugh* and F. M. Ryan, Carnegie Institute of Technology, Pittsburgh, Pennsylvania (Received April 30, 1958).

The magnetic susceptibilities of copper-nickel alloys with up to 39 atomic percent nickel and silver-palladium alloys with up to 10 atomic percent palladium have been measured by the Gouy method at temperatures between 295°K and 2.1°K. Measurements are reported on samples with an impurity content of less than 10 parts per million and on samples with greater impurity content. The susceptibility is nearly temperature independent for all the silver-palladium alloys and for the copper-nickel alloys with up to 27 percent nickel. Higher nickel concentrations give rise to a $1/T$ susceptibility term which is too large to be accounted for by impurity content. The simple band model of Mott adequately describes the susceptibility of the silver-palladium alloys but fails to predict the large temperature-independent paramagnetic contribution of nickel in copper. Exchange energies calculated by the Stoner collective electron method for the copper-nickel alloys, using specific heat and susceptibility data, are somewhat unsatisfactory and suggest the presence of another mechanism.

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INTERACTION OF SLOW ELECTRONS WITH ATOMIC OXYGEN AND ATOMIC NITROGEN. Milton M. Klein, Aerosciences Laboratory, General Electric Company, Philadelphia, Pennsylvania, and Keith A. Brueckner, University of Pennsylvania, Philadelphia, Pennsylvania (Received April 16, 1958).

The polarization potential for an electron in the field of an oxygen atom has been determined by utilizing recent experimental results on the binding energy of O^-. The corresponding polarization potential for atomic nitrogen has been obtained from the results for oxygen by an extrapolation based on the theory of polarization. From these results the photo cross section for O^- and the scattering cross section sections for oxygen and nitrogen have been calculated. The photo cross section is in good agreement with experimental results. The scattering cross sections are compared with other recent calculations.

It is also shown on the basis of the general theory of the photo cross section that the energy dependence of the cross section determines the scattering length for electron scattering. The agreement of the computed photo cross section with experiment therefore is a direct check on the scattering prediction.

A THEOREM FOR GENERALIZED OSCILLATOR STRENGTHS. S. M. Silverman, Geophysics Research Directorate, Air Force Cambridge Research Center (Received April 22, 1958).

A theorem for generalized oscillator strengths is presented. The theorem should prove useful in the interpretation of experimental and theoretical electronic collision cross sections.

VARIATIONAL CALCULATIONS OF THE $2\,^3S$ STATE OF HELIUM. J. Traub and H. M. Foley, Columbia University, New York, New York (Received March 13, 1958).

With a 12-parameter Hylleraas type wave function containing only positive powers, a new calculation has been carried out for the $2\,^3S$ state of helium by the Ritz variational principle. The energy was minimized by a descent process. A nonrelativistic energy of -1.0876088 Hylleraas units was reached as compared with the best previously published value of -1.0876015 Hylleraas units from a 6-parameter function. When mass polarization and α^2ry corrections are included, the 12-parameter function gives an ionization potential 2.52 cm^{-1} less than the experimental value of 38 454.64 cm^{-1}. The electron density at the nucleus is also calculated and compared with the experimental hyperfine spectrum value. All numerical work was carried out on an IBM 650 computer.

ANALYTIC HARTREE-FOCK SOLUTIONS FOR $O^=$. Richard E. Watson, Massachusetts Institute of Technology, Cambridge, Massachusetts (Received April 25, 1958).

$O^=$ ions exist in solids but not in a free state. Analytic Hartree-Fock solutions have been obtained for the ^1S state of $O^=$ which can be applied to work in solids. The solutions utilize stabilizing potential wells of positive charge and the results are compared with other published work.

SPIN OF NEPTUNIUM-238. R. G. Albridge, J. C. Hubbs, and R. Marrus, Radiation Laboratory and Department of physics and Chemistry, University of California, Berkeley, California (Received April 14, 1958).

The atomic-beam magnetic-resonance method has been used to investigate 2.10-day Np^{238} in the low-field or Zeeman region of hyperfine structure. The spin of this nuclide is found to be 2.

PROTON POTENTIAL ANOMALY AND NON-LOCAL POTENTIALS. A. E. S. Green, Los Alamos Scientific Laboratory and The Florida State University, and P. C. Sood, The Florida State University* (Received April 25, 1958).

A detailed study of the variations of the difference between the proton-nuclear potential and the neutron-nuclear potential is carried out using ground state mass data. A previous analysis of mirror nuclei has delimited the problem and suggested that the well difference can neither be constant nor strictly Z dependent. It is found by using analytical techniques as well as calculations performed with the Oak Ridge Oracle that the potential anomaly varies with the symmetry parameter $(N-Z)/A$ and that the proportionality constant does not change very greatly as one proceeds from a static well to a velocity-dependent well. This is somewhat surprising in view of many studies which suggest that the nuclear symmetry energy is considerably influenced by the velocity dependence of the nuclear potential. The relative insensitivity to velocity dependence for nonlocality is attributed to surface effects not taken into account in the analyses of infinite nuclear matter.

A number of possible origins of the well differences are examined including (a) the failure of Koopmans' theorem, (b) the breakdown of electrostatic laws, (c) the presence of Heisenberg forces, (d) the effect of the exclusion principle, and (e) the spin dependence of nuclear forces. It is concluded that the last two effects account for the bulk of the proton potential anomaly. Indeed, these last two effects suggest that the well depths used in neutron scattering and in proton scattering vary individually with the symmetry parameter, and it is suggested that experimentalists attempt to seek out these variations in careful, low-energy scattering experiments.

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PROTON-GAMMA RAY ANGULAR CORRELATIONS IN THE $Si^{28}(d,p\gamma)Si^{29}$ REACTION. H. A. Hill* and J. M. Blair, School of Physics, University of Minnesota, Minneapolis, Minnesota (Received April 30, 1958).

The $Si^{28}(d,p\gamma)Si^{29}$ reaction has been studied by measuring angular distributions of the protons leaving Si^{29} in its first and second excited states and by measuring the correlation between the protons of each group and the gamma rays emitted from the corresponding level. The angular distributions of proton groups p_1 and p_2 exhibit large peaks in the backward direction, indicating that heavy particle stripping may be important. The (p_1,γ) correlations were found to be isotropic to ± 6 percent. This is substantially different from the correlation which has been observed at a higher deuteron energy. The (p_2,γ) correlations are consistent with a 5/2+ level in Si^{29} at 2.03 Mev. One (p_2,γ) correlation shows a 20 percent anisotropy in the plane perpendicular to the direction taken by the recoil nucleus, while another (p_2,γ) correlation is essentially isotropic.

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THE $T(t,\alpha)n,n$ REACTION. Nelson Jarmie and Robert C. Allen, Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received April 25, 1958).

The alpha-particle spectra from the $T(t,\alpha)n,n$ reaction have been measured at a laboratory angle of 30° for triton energies ranging from 0.95 to 2.10 Mev and at laboratory angles of 30°, 60°, 90°, and 120° for 1.9-Mev incident tritons. Absolute cross sections are obtained. Analysis of the spectra in terms of a two-stage process involving the formation and breakup of He^5 is discussed. Evidence for neutron-neutron correlation is observed but no evidence is found for a bound dineutron.

TRITON REACTION CROSS SECTIONS. Robert C. Allen and Nelson Jarmie, Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received April 25, 1958).

Measurements of the $D(t,\alpha)n$ reaction cross section at triton energies of 1.9 and 1.5 Mev have been made to resolve a disagreement among previously published values. Corrections for published values of the $He^4(t,t)He^4$ and

T(*t*, *t*)T scattering cross sections are also suggested.

NEUTRON NONELASTIC CROSS SECTIONS AT 21.0, 25.5, and 29.2 MEV. Malcolm H. MacGregor, William P. Ball,* and Rex Booth, University of California Radiation Laboratory, Livermore, California (Received April 16, 1958).

Neutron inelastic cross sections for eleven elements have been measured in the energy range 21-29 Mev. Corrections were applied to the data by means of a UNIVAC calculation. Since no experimental neutron angular distributions are available in this energy region, optical-model calculations were used for the correction problem. In general the cross sections at 25 Mev are 10% to 20% lower than at 14 Mev.

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SEARCH FOR DELAYED NEUTRONS FROM THE PHOTON BOMBARDMENT OF LITHIUM. G. W. Tautfest, Department of Physics, Purdue University, Lafayette, Indiana (Received April 21, 1958).

A search has been made for delayed neutrons resulting from the reaction $Li^7(\gamma, 2p)H^5$ in an effort to decide whether or not H^5 exists. Delayed neutrons were detected in a moderated counter assembly and scaled in nine delay channels which were variable in width and initial delay. A small signal observed after background subtraction showed no decay within statistics over a period of 0.45 sec. The first 0.05-sec counting period has been analyzed for the presence of a decay of 0.01-sec half-life and compared with a predicted yield of the reaction $Li^7(\gamma, 2p)$ obtained by extrapolation of the measured yields of the reactions $B^{11}(\gamma, 2p)$ and $F^{19}(\gamma, 2p)$. The conclusion is that less than one percent of the expected $Li^7(\gamma, 2p)$ yield can result in a particle-stable final state.

PROTON-PROTON SCATTERING IN THE Bev REGION. G. E. Brown, Institute for Theoretical Physics, University of Copenhagen, Copenhagen, Denmark (Received April 21, 1958)..

Proton-proton scattering in the Bev region is analyzed in terms of an interaction which, at 1 Bev, is taken to be a hard core of radius 0.45 $\times 10^{-13}$ cm together with an external absorption of Gaussian form. The hard core is assumed to disappear with increasing energy and to be replaced by absorption. General features of the data are well reproduced by this simple model.

MASS EFFECT IN *P*-STATE π-*p* SCATTERING. Hong-Yee Chiu, Laboratory of Nuclear Studies, Cornell University, Ithaca, New York (Received February 27, 1958; revised manuscript received May 26, 1958).

The mass difference between charged and neutral mesons has been correlated to a change of the pseudovector coupling constant, giving rise to an additional term in the interaction Hamiltonian. The mass effect in *P*-state π^--*p* scattering is then calculated, using Chew's fixed source theory. A correction to the usual scattering formula for the π^--*p* system is obtained. This correction fails to explain the quantitative feature of the observed difference $(\delta_{13} - \delta_{31})$.

NUCLEON-NUCLEON INTERACTIONS AND POLARIZATION OF HIGH-ENERGY PROTONS ELASTICALLY SCATTERED FROM CARBON. Shoroku Ohnuma, Sloane Physics Laboratory, Yale University, New Haven, Connecticut (Received April 28, 1958).

The transition matrix element in momentum space derived by Riesenfeld and Watson has been used to calculate the polarization and the triple-scattering parameter β of high-energy protons elastically scattered at small angles from carbon. For the nucleon-nucleon phase shifts which represent two-body interactions, those by Signell and Marshak, by Gammel and Thaler, and by Feshbach and Lomon have been considered. In evaluating nuclear as well as Coulomb scattering amplitudes, the first Born approximation has been employed with the assumption that, in carbon, the distribution of protons is equal to that of neutrons. Final results are then independent of the assumed distribution of nucleons. It has been found that while one cannot discriminate between Signell-Marshak and Gammel-Thaler phase shifts, both of them being in semiquantitative agreement with experimental data, Feshbach-Lomon phase shifts may be ruled out because of the wrong sign of the resulting β. Since only the first-order transition matrix element in momentum space has been used in the present work, the calculation does not depend on the optical model potential in the usual sense.

SUBSIDIARY CONDITIONS IN COVARIANT THEORIES. Peter G. Bergmann, Syracuse University, Syracuse, New York, and Allen I. Janis, University of Pittsburgh, Pittsburgh, Pennsylvania (Received April 1, 1958).

We have investigated the effect that coordinate conditions and similar conditions will have on the formal properties of covariant theories. Two distinct types of coordinate conditions were included, those involving first derivatives of the field variables (such as the gauge condition of Lorentz and the coordinate conditions of De Donder) and those algebraic in the field variables (of which the Coulomb gauge is an example).

We have found that with either type of coordinate condition we can construct a variational principle, or a Hamiltonian formalism, which leads to physically meaningful field equations if associated with appropriate initial conditions on a space-like hypersurface. Thus the existence of a properly set Cauchy problem is always assured.

It had been found previously that the infinitesimal invariant transformations of covariant theories form a group and that the coordinate (and similar) transformations represent a normal subgroup. The members of the resulting factor group are in one-to-one correspondence with the true observables of the theory, those dynamical variables which alone possess intrinsic significance without reference to a particular frame of description and whose commutator algebra is presumably reflected in the commutators of the corresponding Hilbert operators of the quantized theory. In this paper we have established the appropriate transformation groups (and their subgroups and factor groups) of a theory with either type of coordinate conditions. We have found that in any of these versions the theory will yield the same observables with the same commutator algebra. One may therefore hope that a quantization scheme based on a theory with subsidiary conditions will be free of the arbitrariness involved in the choice of particular conditions.

SPECTRAL REPRESENTATIONS IN PERTURBATION THEORY I. VERTEX FUNCTION. Robert Karplus, Charles M. Sommerfield, and Eyvind H. Wickmann, Physics Department, University of California, Berkeley, California (Received April 21, 1958).

The vertex operator is examined in lowest order perturbation theory. It is found that, as a function of the invariant momentum transfer $q^2 = q^2 - q_0^2$, it is analytic in a cut plane with the branch point on the negative real axis. A spectral representation (dispersion relation) may therefore be inferred. The threshold of the spectrum depends on the masses of all fields involved unless certain inequalities hold between the masses of the incident and outgoing particles on one hand and the particles in intermediate states on the other; in that case the threshold depends only on the intermediate masses.